The Stimulated Brillouin Scattering during the Interaction of Picosecond Laser Pulses with Moderate Scale-Length Plasmas

by

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1999
A Mama, Isabel y Susana,
Coti y Abuela.
CURRICULUM VITAE

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ABSTRACT

The Stimulated Brillouin Scattering (SBS) instability is studied in moderately short scale-length plasmas. The backscattered and specularly reflected light resulting from the interaction of a pair of high power picosecond duration laser pulses with solid Silicon, Gold and Parylene-N (CH) strip targets was spectrally resolved.

The first, weaker laser pulse forms a short scale-length plasma while the second delayed one interacts with the isothermally expanded, underdense region of the plasma. The pulses are generated by the Table Top Terawatt (TTT) laser operating at 1054 nm (infrared) with intensities up to $5 \times 10^{16}$ W/cm$^2$. Single laser pulses only show Lambertian scattering on the target critical surface. Pairs of pulses with high intensity in the second pulse show an additional backscattered, highly blueshifted feature, associated with SBS. Increasing this second pulse intensity even more leads to the appearance of a third feature, even more blueshifted than the second, resulting from the Brillouin sidescattering of the laser pulse reflected on the critical surface.

The SBS threshold intensities and enhanced reflectivities for P-polarized light are determined for different plasma density scale-lengths. These measurements agree with the convective thresholds predicted by the SBS theory of Liu, Rosenbluth, and White using plasma profiles simulated by the LILAC code.

The spectral position of the Brillouin back- and sidescattered features are determined. The SBS and Doppler shifts are much too small to explain the observed blueshifts. The refractive index shift is of the right magnitude, although more detailed verification is required in the future.
# Table of Contents

**Title** ...................................................................................................................... i  
**Curriculum Vitae** ................................................................................................... iii  
**Acknowledgements** ............................................................................................... iv  
**Abstract** ................................................................................................................. v  
**Table of Contents** .................................................................................................... vi  
**List of Figures** ......................................................................................................... ix  
**List of Tables** ........................................................................................................... xiii

I- **Introduction** ........................................................................................................ 1  
I.1 - Review of Previous Work ..................................................................................... 4  
I.2 - Motivation ............................................................................................................ 5  
I.3 - Conclusions .......................................................................................................... 8  
I.4 - Thesis Outline ...................................................................................................... 9  
References for chapter I ............................................................................................... 11

II- **Theory** ................................................................................................................. 15  
II.1 - Stimulated Brillouin Scattering in a ................................................................. 15  
  A- SBS: Hydrodynamic Model ................................................................................. 15  
  B- The Brillouin Scattered Light Spectrum ............................................................. 24  
  C- The SBS Threshold with Collisional and Landau Damping ......................... 27  
  D- Non-isothermal effects for high-Z plasmas ...................................................... 32  
II.2 - SBS in a Drifting Inhomogeneous Plasma ....................................................... 37  
  A- The Doppler Shifting of the Backscattered Light ............................................. 37  
  B- The Effect of Density, Velocity and Temperature Gradients ......................... 38  
II.3 - Finite-Length Pulse Effects .............................................................................. 44  
  A- SBS Saturation time ........................................................................................... 44  
  B- The Temporal threshold .................................................................................... 52  
II.4 - Seeding the Instability ....................................................................................... 53  
  A- Seeding by the Thermal EM Noise Background ............................................. 53  
  B- Seeding by Accelerated Ion Jets ....................................................................... 54  
  C- The Coupling of SBS to SRS ......................................................................... 56  
II.5 - Translating SBS Theory into an Experiment .................................................. 57  
  A- Ideal and observable spectra .......................................................................... 57  
  B- Measuring the SBS threshold ........................................................................... 60  
References for chapter II .............................................................................................. 63
III- EXPERIMENTAL SET-UP ........................................................................ 66

III.1 - Plasma Formation by Ultrashort Laser pulses ......................... 66
   A- Plasma Formation by Light Absorption .................................. 68
      a- Collisional Absorption ................................................... 68
      b- Resonant Absorption ..................................................... 71
   B- Free Expansion of the Preformed Isothermal Plasma ............... 75
   C- Plasma Heating by the Interaction Pulse .............................. 79
   D- LILAC Simulations ............................................................ 80

III.2 - TTT Chirped Pulse Amplification Laser System ....................... 86

III.3 - Pulse Splitting, Delaying and Stacking ..................................... 87

III.4 - High-Contrast Laser Pulses .................................................... 89

III.5 - The Zeta-Tank ........................................................................ 91

III.6 - Targets .................................................................................. 94

III.7 - Diagnostics .......................................................................... 95
   A- Incident Laser Light ........................................................... 95
      a- Total incident laser energy ......................................... 95
      b- Contrast between preforming and interaction pulses .... 95
   B- Backscattered Light .......................................................... 97
      a- Total backscattered energy ........................................... 97
      b- Time-integrated backscattered spectrum .................... 97
   C- Energy calibration of the Z-tank diagnostics ....................... 98
   D- Density scalelength ............................................................ 101
   E- Laser focus displacement ................................................ 108

References for chapter III ................................................................. 110

IV- RESULTS AND ANALYSIS ................................................................. 113

IV.1 - Silicon Targets ...................................................................... 113
   A- Backscattered Reflectivity ............................................... 113
      a- Single laser pulses ....................................................... 114
      b- Double pulses ............................................................ 115
   B- Spectra ........................................................................... 119
      a- Single laser pulses ....................................................... 119
      b- Double laser pulses with fixed target angle of incidence 121

IV.2 - Gold targets ........................................................................ 133
   A- Backscattered Reflectivities .............................................. 133
   B- Spectra ........................................................................... 135

IV.3 - Parylene-N (CH) .................................................................. 139
   A- Backscattered Reflectivities .............................................. 139
   B- Spectra ........................................................................... 141

IV.4 - Common Features of the Experimental Data ......................... 145
   A- SBS Intensity Thresholds ................................................ 145
   B- Spectral Features ............................................................. 149

References for chapter IV ............................................................... 153
V. SUMMARY AND CONCLUSIONS ............................................. 154
  V.1 - Summary ......................................................... 154
  V.2 - The Brillouin Thresholds ................................. 154
  V.3 - Brillouin Spectra ........................................... 156
  V.4 - Conclusions .................................................. 157
References for chapter V ............................................. 159

APPENDIX
A- MEASUREMENTS OF BACKSCATTERED LIGHT NEAR 351 NM IN
OMEGA LONG SCALELENGTH PLASMA EXPERIMENTS ................. 160
  Abstract .......................................................... 160
  A.1 - Introduction .................................................. 161
  A.2 - Experimental conditions ................................... 162
    A- Long Scale Length Plasma Formation ......................... 162
    B- Interaction Beam ............................................ 163
    C- Experimental set-up ....................................... 163
  A.3 - Characteristics of the Backscattered Light ............... 164
  A.4 - Analysis and discussion .................................... 168
    A- Total backscattered energy and reflectivity .............. 172
    B- Thomson ion feature peak .................................. 172
    C- SBS Threshold ............................................... 174
  A.5 - Conclusions ................................................ 177
    Acknowledgments ................................................ 177
References for Appendix A ........................................... 178
List of Figures

Fig. I.1 SBS backscattering on a ion-acoustic density fluctuation. 3
Fig. I.2 Schematic of the short laser pulse experiment for the measurement of Brillouin backscattered light. 7

Fig. II.1 A wavenumber vector diagram of SBS. 23
Fig. II.2a,b Quartic and strong-field solutions to the Brillouin dispersion relation. 25
Fig. II.2c Quartic and strong-field solutions to the Brillouin dispersion relation. 26
Fig. II.3 Dissipative SBS threshold intensity for a homogeneous stationary fully-ionized Si plasma for normally incident infrared light ($\lambda_L = 1.054 \, \mu\text{m}$). 31
Fig. II.4 Ratio of the square of the SBS growth rates of the non-local transport model to the isothermal model. 36
Fig. II.5 Region of resonance for SBS in an inhomogeneous drifting plasma. 40
Fig. II.6 Convective SBS threshold in an inhomogeneous, fully ionized Silicon plasma. 43
Fig. II.7 SBS response to a short duration laser pulse incident on a finite homogeneous plasma. 45
Fig. II.8 SBS response to a long duration laser pulse incident on a finite homogeneous plasma. 47
Fig. II.9 SBS impulse response of a finite scalelength, homogeneous Silicon plasma to a seminfinite laser pulse. 50
Fig. II.10 Temporal SBS threshold intensity for a homogeneous stationary fully-ionized Si plasma for normally incident 1.6 ps FWHM laser pulse. 52
Fig. II.11 Brillouin scattering from an ion beam decaying into ion-acoustic waves through an streaming instability. 55
Fig. II.12 Observable stimulated Brillouin light scattering by a plasma expanding from a plane solid target. 58
Fig. II.13 Appearance of the observable SBS spectrum. 59
Fig. II.14 The ideal behavior of the total SBS signal with increasing laser interaction intensity. 61
Fig. II.15 Multiple SBS growth curves appear on the actual measurement of the backscattered intensity. 62

Fig. III.1 Plasma formation and SBS interaction sequence. 67
Fig. III.2 Absorption of a obliquely incident laser beam in an inhomogeneous plasma. 69
Fig. III.3 Collisional fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a 1 keV Silicon plasma. 70
Fig. III.4 Resonant fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a plasma. 72
Fig. III.5a,b Resonant combined with collisional scattered (non-absorbed) fraction at different electron density scalelengths for a P-polarized 1054 nm laser beam interacting with a Silicon plasma. 73
Fig. III.5c Resonant combined with collisional scattered (non-absorbed) fraction at different electron density scalelengths for a P-
polarized 1054 nm laser beam interacting with a Silicon plasma.

Fig. III.6- Ion density ratio ($n_i/n_e$) and Mach expansion speed for planar isothermal expansion of a plasma. 74

Fig. III.7 Ion density ratio ($n_i/n_e$) and Mach expansion speed for spherical isothermal expansion of a plasma. 77

Fig. III.8a LILAC calculated $n_i/n_{cr}$ profile. 81
Fig. III.8b LILAC calculated plasma expansion speed profile. 81
Fig. III.8c LILAC calculated electron temperature profile. 82
Fig. III.8d LILAC calculated ion temperature profile. 82
Fig. III.9a LILAC calculated $n_0/n_{cr}$ profiles. 84
Fig. III.9b LILAC calculated plasma expansion velocity profiles. 84
Fig. III.9c LILAC calculated electron temperature profiles. 85
Fig. III.9d LILAC calculated ion temperature profiles. 85
Fig. III.9e LILAC calculated plasma expansion velocity profiles. 86
Fig. III.10 The TTT Chirped Pulse Amplification laser system. 87
Fig. III.11 Pulse splitting, delaying and stacking set-up. 89
Fig. III.12 Saturable absorber transmission curve for the laser beam. 90
Fig. III.13 TTT laser pulse before (low contrast) and after (high contrast) the passage through the saturable absorber. 91

Fig. III.14 Diagnostics set-up to measure both incident and backscattered light at TTT Zeta tank. 92
Fig. III.15 Normal transverse image of the laser focus (far field). 93
Fig. III.16a Transient-Stillcope trace for a single laser pulse. 96
Fig. III.16b Transient-Stillcope trace for a double laser pulse, separated by 530 ps. 96

Fig. III.17 Nd-YLF Oscillator spectra for different slit apertures. 98
Fig. III.18 Calibration set-up for all energy diagnostics at the Z-tank. 99
Fig. III.19a,b Calibration plot for total and spectral backscattered energy. 100
Fig. III.19c Calibration plot for total and spectral backscattered energy. 101
Fig. III.20 Integrating Sphere set-up to measure density scalelengths $L_n$ by maximum resonant absorption. 102

Fig. III.21 Integrating Sphere measured transmittances for a single laser pulse ($\Delta \tau = 0$ ps) at different Silicon target angles. 104
Fig. III.22 Integrating Sphere measured transmittances for double laser pulses ($\Delta \tau = 1260$ ps) for different Silicon target angles. 104
Fig. III.23 Measured and LILAC-calculated electron density scalelengths $L_n$ at different delays $\Delta \tau$ for Silicon planar targets. 105
Fig. III.24 The expansion of a plasma near a solid target surface. 106
Fig. III.25 Electron temperature $T_e$, measured near the critical surface and LILAC-calculated average, at different delays $\Delta \tau$ for Silicon planar targets. 107

Fig. III.26 CCD images of the incident laser focus reflected on a Silicon planar target for a single pulse and for a double pulse separated by $\Delta \tau = 1260$ ps. 108

Fig. IV.1 Total backscattered intensity—PIN diode measured—for single laser pulses interacting with a planar SI target. 114
Fig. IV.2a Total reflectivity for double laser pulses separated by $\Delta \tau = 260$ ps interacting with a planar SI target. 116
Fig. IV.2b Total reflectivity for double laser pulses separated by $\Delta \tau = 510$ ps interacting with a planar SI target. 117
Fig. IV.2c Total reflectivity for double laser pulses separated by $\Delta \tau = 760$ ps interacting with a planar SI target. 117
Fig. IV.2b,c Total reflectivity for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar SI target. 118
Fig. IV.2d,e Total reflectivity for double laser pulses separated by $\Delta \tau = 1260$ ps interacting with a planar SI target. 118
Fig. IV.3 Backscattered spectrum (averaged) for single laser pulses interacting with a planar Si target. 120
Fig. IV.4a Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Si target. 122
Fig. IV.4b Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Si target. 122
Fig. IV.4c Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Si target. 123
Fig. IV.4d Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Si target. 124
Fig. IV.4e Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Si target. 125
Fig. IV.5a,b Backscattered spectrum for double laser pulses separated by \( \Delta t = 260 \) ps interacting with a planar Si target. 126
Fig. IV.5c,d Backscattered spectrum for double laser pulses separated by \( \Delta t = 260 \) ps interacting with a planar Si target. 127
Fig. IV.6a,b Backscattered spectrum for double laser pulses separated by \( \Delta t = 510 \) ps interacting with a planar Si target. 128
Fig. IV.7a,b Backscattered spectrum for double laser pulses separated by \( \Delta t = 760 \) ps interacting with a planar Si target. 129
Fig. IV.7c,d Backscattered spectrum for double laser pulses separated by \( \Delta t = 760 \) ps interacting with a planar Si target. 130
Fig. IV.8a Backscattered spectrum for double laser pulses separated by \( \Delta t = 1260 \) ps interacting with a planar Si target. 131
Fig. IV.8b,c Backscattered spectrum for double laser pulses separated by \( \Delta t = 1260 \) ps interacting with a planar Si target. 132
Fig. IV.8d Backscattered spectrum for double laser pulses separated by \( \Delta t = 1260 \) ps interacting with a planar Si target. 133
Fig. IV.9 Backscattered intensity for single laser pulses interacting with a planar Au target. 134
Fig. IV.10 Backscattered reflectivity for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Au target. 135
Fig. IV.11 Backscattered spectrum for single laser pulses interacting with a planar Au target. 136
Fig. IV.12a Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Au target. 137
Fig. IV.12b,c Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar Au target. 138
Fig. IV.13 Backscattered intensity for single laser pulses interacting with a planar CH target. 140
Fig. IV.14 Backscattered reflectivity for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar CH target. 141
Fig. IV.15 Backscattered spectrum for single laser pulses interacting with a planar CH target. 142
Fig. IV.16a,b Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar CH target. 143
Fig. IV.16c,d Backscattered spectrum for double laser pulses separated by \( \Delta t = 1010 \) ps interacting with a planar CH target. 144
Fig. IV.17 LILAC-calculated LRW SBS thresholds for a Silicon plasma formed at \( \Delta t = 1010 \) ps before Interaction. 146
Fig. IV.18 Measured and calculated SBS thresholds for Silicon at different interval between laser pulses. 147
Fig. IV.19 Measured and calculated (LRW model) SBS thresholds for Silicon at different electron density scalelengths. 148
Fig. IV.20 Measured and calculated absolute spectral shifts of Brillouin backscattered light for Silicon. 151
Fig. IV.21 Measured spectral shifts of Brillouin back- and sidescattered light for Silicon at \( \Delta t = 1010 \) ps. 152
<table>
<thead>
<tr>
<th>Fig. IV.22</th>
<th>Spectral shift of Brillouin backscattered and sidescattered features for Silicon.</th>
<th>152</th>
</tr>
</thead>
<tbody>
<tr>
<td>Fig. A.1</td>
<td>Diagnostics set-up for Omega Stimulated Brillouin backscattering experiments (E6696).</td>
<td>164</td>
</tr>
<tr>
<td>Fig. A.2a</td>
<td>Backscattered spectrum for OMEGA shot #23978 (E6639).</td>
<td>165</td>
</tr>
<tr>
<td>Fig. A.2b</td>
<td>Backscattered spectrum for OMEGA shot #24151 (E6638).</td>
<td>166</td>
</tr>
<tr>
<td>Fig. A.2c</td>
<td>Backscattered spectrum for OMEGA shot #24258 (E6637).</td>
<td>167</td>
</tr>
<tr>
<td>Fig. A.3</td>
<td>Backscattered spectrum for OMEGA shot #24110 for the interaction with a Solid alignment target (E6636).</td>
<td>169</td>
</tr>
<tr>
<td>Fig. A.4</td>
<td>Fractional backreflected energy (E6601).</td>
<td>170</td>
</tr>
<tr>
<td>Fig. A.5</td>
<td>Temporal profile of OMEGA shots 24665, 24668.</td>
<td>171</td>
</tr>
<tr>
<td>Fig. A.6</td>
<td>Thermal Thomson scattering ion peak energies (E6641).</td>
<td>173</td>
</tr>
<tr>
<td>Fig. A.7</td>
<td>Calculated SAGE plasma density and expansion velocity profiles (E6640).</td>
<td>176</td>
</tr>
</tbody>
</table>
List of Tables

Table III.1  LILAC calculated plus integration Sphere measured Scalelengths near the plasma critical surface. ........................................... 107

Table IV.1  Measured and calculated SBS thresholds for Silicon, Gold, and Parylene-N (CH) at $\Delta t = 1010$ ps. .............................................. 149
I- INTRODUCTION

STIMULATED BRILLOUIN SCATTERING (SBS) is a parametric instability that can occur in the subcritical region of a plasma irradiated by an electromagnetic (EM) wave. A transverse electromagnetic wave scatters parametrically from a longitudinal ion-acoustic (IA) wave; thus, this process may be regarded as the low-frequency counterpart to the Stimulated Raman Scattering (SRS), in which EM waves are scattered parametrically by electron plasma (EP) waves. The SBS instability is important in the progress of laser fusion. By scattering the laser light away from the absorption region near the critical density (where the laser light frequency $\omega_0$ is equal to the electron plasma frequency $\omega_{pe}$) it may reduce the amount of laser energy available for coupling to the target.

The Stimulated Brillouin Scattering (SBS) instability can be a serious energy loss mechanism in laser-driven inertial confinement fusion, both in the direct-drive and in the indirect-drive configurations. Under ignition conditions, using the direct drive configuration, the pellet’s expanding coronal plasma has a long-scalelength ($L_\perp \approx 10^4 \lambda_L$). When using the indirect-drive configuration, similar or even longer scalelengths appear in the expanding plasma coming from the internal walls of the Hohlraum. These plasmas are created by several nanosecond duration laser pulses and the SBS instability reaches saturation for most of the duration of the interaction pulse. Therefore, an important body of theoretical and experimental work—as described in section I.1 and its references—has been performed to understand the behavior of this instability.

Ordinary Brillouin Scattering arises as a result of the beating of the incident EM wave with a low-frequency ion-acoustic mode such that frequency and wavenumber matching Manley-Rowe relations:

$$\omega_0 = \omega_{ia} + \omega_s, \quad k_0 = k_{ia} + k_s.$$  \hspace{1cm} (I.1)

are satisfied. Here, $(k_0, \omega_0)$ are the incident EM wave wavenumber and frequency, $(k_{ia}, \omega_{ia})$ are the ion-acoustic wave wavenumber and frequency, and $(k_s, \omega_s)$ are the scattered EM wave wavenumber and frequency. These quantities obey the following dispersion relationships

$$\omega_0^2 = k_0^2 c_s^2 + \omega_{pe}^2, \quad \omega_{ia} = k_{ia} c_s.$$
\[ \omega_s^2 = k_s^2 c^2 + \omega_{pe}^2. \] (I.2)

The electron plasma frequency is defined as

\[ \omega_{pe}^2 = \frac{4 \pi n_e e^2}{m_e}, \] (I.3)

where \( m_e, n_e \) and \( e \) are respectively the mass, number density and charge of the electron. Electromagnetic waves cannot propagate where the electron density is over critical. The plasma critical surface is located where the condition

\[ \omega_{pe}^2 = \omega_0^2 \]

is satisfied; therefore the critical electron density is

\[ n_{cr} = \frac{m_e}{4 \pi e^2} \omega_0^2 = 1.12 \cdot 10^{21} / \lambda_{L[\mu m]}^2 [\text{cm}^{-3}]. \] (I.4)

The ion-sound speed is

\[ c_s = \sqrt{\frac{Z_i kT_e + kT_i}{m_i}}. \] (I.5)

where \( T_e, T_i \) are respectively the electron and ion temperature (\( k \) is the Boltzmann gas constant), \( m_i \) is the ion mass, and \( Z_i \) is the ion charge state.

Figure I.1 shows the development of the Stimulated Brillouin Scattering instability schematically. A fraction of the incident laser light scatters—Bragg-like—from the ion density fluctuations created by ion-acoustic noise present in the plasma, creating a small amplitude scattered EM wave, with all waves obeying the relationships I.1 and I.2. This process is plain (non-stimulated) Brillouin scattering. Raising the intensity of the incident laser light, the Brillouin scattering process becomes unstable (stimulated) when the ponderomotive force,\(^5\)

\[ F_{pond} = \frac{e^2}{4m_e \omega_{pa}^2} \nabla \langle E_{L,sc}^2 (x) \rangle. \] (I.6)

is strong enough to increase the density fluctuations of the original ion-acoustic mode. This force results from the spatially varying superposition of the incident pump EM wave and the scattered EM wave electric field denoted \( E_{L,sc}(x) \). Beating at the frequency \( \omega_{pa} \) and wavenumber \( k_{pa} \) of the ion-acoustic wave, the ponderomotive force increases the amplitude.
of the ion density fluctuations. Higher ion density fluctuations provide a stronger scattering of the incident EM wave. A feedback mechanism is established and scattered EM wave and ion-acoustic wave grow exponentially together until some saturation level is reached.

Characteristic features of SBS include:

- The frequency shift $\Delta \omega$ of the scattered EM wave is small because the ion-acoustic frequency $\omega_{ia}$ is much smaller than the frequency $\omega_L$ of an EM wave when the two wavelengths are comparable.

![Diagram of SBS backscattering](image)

**Fig. I.1 - SBS backscattering on a ion-acoustic density fluctuation.**
A- Incident laser light interacts with an ion-acoustic wave in the plasma, resulting in the Bragg-like scattering of part of this EM wave. B- The electric field resulting from the superposition of the incident and backscattered light electric field generates a ponderomotive pressure $\nabla \langle E^2 \rangle$. C- This force pushes the plasma out of the regions of maximum $\langle E^2 \rangle$. This feedback effect enhances the original ion-acoustic density fluctuation and increases the reflectivity of the Bragg grating.

- The small frequency shift $\Delta \omega$ in the scattered EM wave means that the energy transferred to the ion acoustic wave is small.
• The SBS instability can potentially reflect a very large fraction of the incoming EM radiation.

• SBS is easier to excite than other parametric instabilities because ion-acoustic waves are low frequency and they have a smaller excitation energy than electron-plasma or electromagnetic waves.

• SBS can grow in the entire subcritical density region (where \( \omega_0 > \omega_{pe} \)) of the plasma.

1.1 - Review of Previous Work

The Stimulated Brillouin Scattering (SBS) instability has been extensively studied due to its potential as a serious energy loss mechanisms in laser driven inertial confinement fusion. Under direct-drive ignition conditions, the plasma has a long-scalelength (\( L_n \sim 10^4 \lambda_L \)) and is created by several nanosecond duration laser pulses. The SBS instability can reach saturation for most of the duration of the interaction pulse. Extensive theoretical \(^6.7.8.9.10\) and experimental work has studied the appearance of SBS in the infrared (\( \lambda_L = 1054 \text{ nm} \)), \(^11.12.13.14.15.16\) green (\( \lambda_L = 531 \text{ nm} \)) \(^17\) and ultraviolet wavelengths (\( \lambda_L = 351 \text{ nm} \)). \(^18.19.20.21.22.23.24\) The conditions in various SBS experiments have varied widely with respect to both the incident EM wave (frequency, pulse length, power, focal spot diameter, laser optics f-number) and the targets (material, geometry, condition). The observed saturated reflectivities range from a few percent to more than 50%. Thresholds, seeding mechanisms and backscattered fractions seem to be uncorrelated with almost any relevant experimental parameter, and there is not much agreement between the observed SBS data and various theoretical predictions over a large range of plasma and laser conditions. \(^25\)

Ultrashort pulse laser-plasma interactions for SBS experiments have the advantage of shortening the growth time of the instability and hence avoiding saturation. The instability is limited to the linear growth regime and so its results should correlate better to the simplest theoretical models of SBS. However, for short scalelength plasmas and ultrashort laser pulses, the theoretical and experimental body of work in SBS is less comprehensive than the thorough theoretical and experimental studies of resonant absorption and reflectance phenomena \(^26.27.28.29.30.31.32.33.34.35.36.37.38.39.40.41\) in the interaction of ultrashort pulses with short scalelength plasmas.
CHAPTER I - INTRODUCTION

Two relevant short interaction pulse, short scalelength plasma, SBS experiments have been performed—the first one in 1993 by Baldis et al. and the second one in 1994 by Baton et al.

In Baldis experiment—at $\lambda_L = 1054$ nm—two laser pulses ($10^{13}$ W/cm$^2 \leq I_L \leq 2 \cdot 10^{15}$ W/cm$^2$) were used. The first one was a long (3 ns) plasma preforming pulse and the second one was a short (8-10 ps) interaction pulse, which was delayed by 10 ns. Plasmas with $n_e < 0.1 \, n_{cr}$ and $T_e = 300$ eV were produced from the ablation of thin plastic (CH) foils. The reflectivities at the higher intensity levels ($< 10\%$ at $10^{15}$ W/cm$^2$) were lower than the theoretical predictions; an effect explained by nonlinear saturation processes in the plasma not incorporated in the model. On the contrary, the reflectivities at low intensities ($\sim 10^{-4}\%$ at $10^{14}$ W/cm$^2$) were higher than those predicted. The parametric decay of coupled plasma waves from the larger growth rate SRS to ion-acoustic ones, and/or localized plasma heating by the interaction beam may lead to an enhanced ion-acoustic noise background level from which SBS could grow.

Baton carried out a shorter pulse length experiment to test the SRS–SBS coupling hypothesis. A 1.2 ps interaction pulse ($5 \cdot 10^{14}$ W/cm$^2 \leq I_L \leq 2 \cdot 10^{17}$ W/cm$^2$) was used in a preformed plasma ($n_e = 0.3–0.5 \, n_{cr}$ and $T_e = 100$ eV, $T_i = 60$ eV) created by exploding thin plastic (CH) foils. The results are similar to those of Baldis et al., with lower than predicted, saturated reflectivities at high intensities ($\sim 10\%$ at $10^{15}$ to $10^{16}$ W/cm$^2$) and higher than predicted reflectivities at low intensities ($\sim 10^{-4}\%$ at $10^{14}$ W/cm$^2$). These results were explained by the seeding of the SBS instability by non-thermal noise with a laser intensity dependence $I^{\alpha}$ ($1 < \alpha < 2$), probably originated in SBS/SRS coupling.

1.2 - MOTIVATION

This thesis reports on a series of experiments have been performed on the short pulse Table Top Terawatt (TTT) laser to study the laser intensity threshold and spectral characteristics of the Stimulated Brillouin Scattering (SBS) instability and its dependence on the plasma’s density scalelength $L_n$, and target composition. The experiments were performed at 1054 nm with intensities up to $5 \cdot 10^{16}$ W/cm$^2$. The backscattered and specularly reflected light resulting from the interaction of two high power picosecond duration laser pulses with solid Si, Au and CH planar targets was measured spectrally. The first pulse created a preformed plasma of density scalelength 30-300 laser wavelengths with which the second pulse interacted.
An ultrashort (1 ps) preforming laser pulse created a short scale length plasma. After a variable time delay (0-1300 ps) an ultrashort (1 ps) interaction laser pulse excited the SBS instability in the isothermally expanded plasma. A critical surface was present in the plasma, which allowed a study of the behavior of SBS for near critical density. The SBS reflectivity and spectral features were measured for several laser intensities, time delays and contrasts between preforming and interaction laser pulses and polarizations.

Creating the plasma conditions needed for the appearance of the SBS instability using only a single intense ultrashort laser pulse requires extremely high intensities. During an ultra-short (~1 ps), intense laser pulse (>10\(^{15}\) W/cm\(^2\)) interacting with a solid target, the plasma is initially formed at densities higher than critical, and during the laser pulse duration its density scalelength is significantly smaller than the laser wavelength (\(L_n/\lambda_L \sim 0.1\)).\(^{49,50}\) The subcritical region is too small (\(L_n/\lambda_{IA} \sim 0.2\)) to allow the formation of an ion-acoustic wave satisfying the wavenumber Manley-Rowe condition (Eqn. I.1b). With these plasma conditions the SBS instability cannot to be excited, amplified and propagated. An underdense plasma scalelength of several laser wavelengths is needed for the onset of SBS, for both enlarging the subcritical density scalelength size to several laser wavelengths and to lower the laser threshold intensity to attainable levels, as the convective threshold scales as \(I_{SBS} \propto 1/L_n\).

If, after the plasma formation, the second intense laser pulse is delayed for a long enough time (> 100 ps) the plasma is predicted to expand to a longer (10-100\(\cdot\)\(\lambda_L\)) density scalelength and SBS will be excited at lower intensities. The use of two ultrashort laser pulses—the first one for plasma formation (of intensity \(I_{for}\)), and the second one (of intensity \(I_{int}\)) for interaction—separated in time offers a set of advantages:

- As the delay between pulses \(\Delta \tau\) is varied over a large range (0-1500 ps), the corresponding plasma scalelength \(L_n\) \((\propto \Delta \tau)\) at interaction time can be controlled.

- The contrast between preforming \(I_{for}\) and interaction \(I_{int}\) laser pulses can be varied (0.01 < \(I_{for}/I_{int}<100\)), and with it the density profile of the preformed plasma and the interaction intensity of the second laser pulse. Normally, \(I_{for}\) is kept constant—to preserve the initial plasma parameters—while \(I_{int}\) is varied.

The control of the delay and contrast between the laser pulses allowed a determination of the relationship between the laser intensity threshold for the SBS instability and the plasma's measured scalelength, and calculated ion and electron temperatures. Also, the seeding intensity level of thermal noise associated to the random ion-acoustic fluctuations
of the plasma from which the SBS instability grows is dependant on its initial electron and ion temperatures (see eqns. II.4.2-3 and fig. I.1).

![Schematic diagram of laser pulse and plasma interaction](image)

**Fig. I.2 - Schematic of the short laser pulse experiment for the measurement of Brillouin backscattered light.**

The figure shows a schematic of the experiment. Single laser pulses originate in a Nd-YLF oscillator. They are split in two separate pulses, delayed one with respect to the other and stacked by a Michelson-like interferometer set-up. They are amplified by Chirped Pulse Amplification. The pedestal superimposed to the amplified pulses is suppressed by the use of a saturable absorber cell. Finally the two laser pulses hit the planar target. The first pulse preforms a plasma. The delayed second pulse scatters in the expanded plasma. The scattered light then is measured spectrally and energetically.

This twin-pulse experimental configuration is shown schematically in fig. I.2. It has some distinctive advantages over previous ones:

- The short preforming pulse creates an initially hot electron ($T_e = 1$ keV), cool ion ($T_i = 1$ eV), very overdense ($n_e = 10^3 \cdot n_{cr}$) plasma, which cools quickly and expands isothermally ($T_e = 50$ eV, $T_i = 20$ eV), with a monotonically decreasing density profile of scalelength $L_n \propto \Delta \tau$.

- The short laser interaction pulse duration minimizes (but does not eliminate) the hydrodynamic and thermal evolution of the plasma during the period while SBS takes place.
• The interaction laser pulse, although ultrashort (~1.6 ps), is long enough to allow full illumination of the expanded preformed plasma, without hindering the convective saturation of the Brillouin instability (see §II.3A).

• An expanded, underdense plasma increases the collisionality ($\propto n_e^3/T_e^{32} = 10^3$ times that at plasma formation—see eqn. II.1.45) of its component species, due to a cooler electron population. However, the collisional damping effects in both the ion-acoustic and scattered electromagnetic waves are small when compared to the convective effects due to the plasma expansion, and this allows the dissipative intensity threshold to be lower (see § II.1C) than the corresponding convective SBS threshold.

### 1.3 - Conclusions

SBS experiments using two ultrashort (~1.6 ps) laser pulses separated by a time interval between 100 and 1300 ps have some definite advantages over previous experiments using single or double long pulses (FWHM $> 10$ ns), as discussed in § I.2. Cleaner and more predictable density, temperature and expansion profiles can be obtained by the plasma formation pulse, while the second interaction pulse keeps the plasma illuminated long enough to reach convective saturation (see § II.3A). Some general conclusions can be reached:

• If the the thermal profile of a plasma is well known at interaction time with a short laser pulse, the Liu-Rosenbluth-White (LRW) isothermal model is able to predict the convective SBS intensity threshold.

Preforming a plasma with another short laser pulse provides a clean, quasi-isothermal plasma, of predictable characteristics. However, the use of the proper diagnostics to find not only the density, but also the expansion velocity and temperature profile of the plasma is needed to provide a better SBS threshold value prediction using the LRW isothermal model.

• The Doppler expansion—even at highly supersonic speeds—of the Brillouin resonant regions of the plasma cannot explain the high blueshifting of SBS features usually found in Brillouin backscattering experiments. The additional high level of blueshifting present on the measured spectra could be explained by the change of
the plasma's refraction index during the time while the interaction laser pulse moves in and out of the SBS resonance region.

Even very approximate hydrodynamic models give a very reasonable prediction to the refraction index shift; although the use of more detailed, but computationally taxing kinetic models,\textsuperscript{51,52} coupled to a better knowledge of the plasma profile could provide a better description of the Brillouin scattered light spectrum.

- The measurement of Brillouin growth rates was very difficult due to the problems of controlling and reproducing the plasma profile consistently at the interaction time.

As a result is impossible to isolate and identify sets of data which have the same initial plasma profile at interaction time, so that only the interaction intensity could be the only variable in play when trying to find the corresponding growth rate. Very stable and reproducible laser pulse characteristics—as focus intensity profile and pulse peak intensity—, plus good quality target surface uniformity and homogeneity are preconditions to achieve a successful measurement of the SBS growth rates.

\section*{I.4 - Thesis Outline}

In Chapter II, the theory of SBS in both homogeneous and inhomogeneous plasmas, is analyzed. General and particular expressions for SBS resonant modes, growth rates, and instability thresholds are found for both cases. The consequences of collisions, Landau damping, ion-acoustic seeding sources, isotopic composition, plus finite laser pulse-length effects are discussed. Using these results, a conceptual explanation of the attainable SBS measurements is shown for both ideal and real experimental conditions.

In Chapter III, the methods and equipment used to perform the experiment are presented. The formation of a plasma by ultrashort laser pulses with solid planar targets plus the subsequent isothermal expansion are analyzed theoretically, numerically and experimentally. The TableTop Terawatt (TTT) Chirped Pulsed Amplification laser system is described. The pulse splitting and stacking using a Michelson interferometer is explained. The generation of high contrast laser pulses by the use of a saturable absorber is described. The Zeta vacuum chamber and its targets are described. The spectral and energy diagnostics for both incoming and backscattered light are described, including their
calibration. The measurement of the preformed plasma scalelength from resonance absorption through the use of an integrating sphere is presented.

In Chapter IV, the SBS spectral and energy measurements are shown and their features classified and compared with the SBS theory presented in the previous two chapters.

In Chapter V, a summary of the main features of the experimental data is presented and an analysis of the correlations between experimental results and theoretical predictions presented in the previous chapters is performed. The validity of the theoretical models and their features are discussed in terms of the observed SBS data.

In Appendix A, the results of a SBS experiment done on the 24-beam Omega laser are presented. This experiment was performed with long scalelength plasmas (>1000 \(\lambda_L\)) and with UV (\(\lambda_L = 351\) nm), long duration (~600 ns) laser pulses.
REFERENCES FOR CHAPTER I


4 W. L. Kruer, ibid., ch. 8, p. 87.

5 W. L. Kruer, ibid., ch. 6, pp. 58-62.


II- Theory

In this chapter, the theory of SBS in both homogeneous and inhomogeneous plasmas is analyzed. General and particular expressions for SBS resonant modes, growth rates, and instability thresholds are found for both cases. The consequences of collisions, Landau damping, ion-acoustic seeding sources, plasma isotopic composition and ionization state plus the finite laser pulse-length effects are included. Using these results, a qualitative description of the attainable SBS measurements is shown for both ideal and real experimental conditions.

II.1 - Stimulated Brillouin Scattering in a Stationary Homogeneous Plasma

A- SBS: Hydrodynamic Model

In the Stimulated Brillouin Scattering (SBS) instability, incident electromagnetic waves \((\omega_1, \mathbf{k}_1)\) interact inelastically with the low frequency ion-acoustic waves \((\omega_{ia}, \mathbf{k}_{ia})\), resulting in scattered electromagnetic waves \((\omega_s, \mathbf{k}_s)\)—which have slightly lower frequencies than those of the incident electromagnetic waves.

To describe SBS, a hydrodynamic model similar to that described in Krueker1,2 is presented. This model is generalized to take account of the effects on SBS of a plasma composed of several ion species of different isotopic composition and ionization state. More detailed, but computationally more taxing, kinetic models for multispecies SBS, can be found in refs. 3, 4. These plasmas are of importance to ICF, because plastic (CHx) and, occasionally, glass (SiO2) pellets are used to contain the fusion fuel. A multispecies plasma is formed in the expanding corona which implodes the fuel to ignition.

SBS must satisfy the Manley-Rowe relations—conservation of momentum and energy—matching the wavenumber and frequency of the incident light wave to those of the scattered light and ion acoustic waves.

\[
\omega_1 = \omega_s + \omega_{ia},
\]

\[
\mathbf{k}_1 = \mathbf{k}_s + \mathbf{k}_{ia}.
\]

Equation (I.1)
Consider a plasma formed by a collection of $N$ different species of ions (each of ion number density $n_{i0}$ and charge state $Z_i$, $1 \leq i \leq N$). The unperturbed plasma is neutral (net electric charge $\sum_{1 \leq i \leq N} Z_i n_{i0} - n_{e0} e = 0$), warm—for electron and all ion species ($T_{i,e} = 0$, $1 \leq i \leq N$)—, stationary, homogeneous and uniform. A large amplitude, coherent monochromatic EM wave illuminates it.

The coupling of the incoming EM wave (vector potential $A_1$) to the scattered EM wave (vector potential $\tilde{A}$) by the electron number density fluctuation $\tilde{n}_e$ is given by\(^{5,6}\)

$$\left( \frac{\partial^2}{\partial t^2} - c^2 \nabla^2 + \omega_{pe}^2 \right) \tilde{A} = -\frac{4\pi e^2}{m_e} \tilde{n}_e A_1.$$  \(\text{(II.1.1)}\)

Only the electron response is considered as the $i$-th ion species ($1 \leq i \leq N$) response to the high frequency EM field is smaller by $Z_i m_e/m_i$.

For the Brillouin instability, the electron density fluctuation is a low frequency one, associated with the ion-acoustic wave. The momentum equation for the electrons is

$$\frac{\partial u_e}{\partial t} + u_e \cdot \nabla u_e = -\frac{e}{m_e} \left[ E + \frac{u_e \times B}{c} \right] - \frac{\nabla p_e}{n_{e0} m_e},$$

where $u_e$ is the electron fluid velocity, $E (= \nabla \phi - \partial A/\partial t)$, $B (= \nabla \times A)$ are the total electric and magnetic fields felt by the electrons, and $p_e = p_{e0} + \tilde{p}_e$ is the electron pressure. Separating the electron fluid velocity $u_e$ into longitudinal ($u_L$) and transverse ($eA/mc$) components, neglecting the electron inertia ($\partial u_L/\partial t \rightarrow 0$, due to the low frequency of the oscillation), considering the electron fluid to be isothermal ($\tilde{p}_e = \tilde{n}_e T_e$), letting $n_e = n_{e0} + \tilde{n}_e$, $A = A_L + \tilde{A}$, $\phi = \tilde{\phi}$, and taking the divergence of the result, gives the linearized electron momentum equation\(^7\)

$$\nabla^2 \tilde{\phi} - \frac{T_e}{en_{e0}} \nabla^2 \tilde{n}_e = \frac{e}{m_e c^2} \nabla^2 (A_L \cdot \tilde{A}).$$  \(\text{(II.1.2)}\)

The plasma is composed of $N$ different ion species, of total ion number density $n_i = n_{i0} + \tilde{n}_i$ ($1 \leq i \leq N$). Poisson’s equation gives

$$\nabla^2 \tilde{\phi} = -4\pi e \left( \sum_{1 \leq i \leq N} Z_i \tilde{n}_i - \tilde{n}_e \right).$$  \(\text{(II.1.3)}\)

The equations of state for all ion species are isothermal ($\tilde{p}_i = \tilde{n}_i T_i$). The linearized mass and momentum conservation equations for the $i$-th ion fluid ($1 \leq i \leq N$), are
\[ \frac{\partial}{\partial t} \tilde{n}_i + n_{i0} \nabla \cdot \tilde{u}_i = 0, \]  

(II.1.4)

\[ \frac{\partial}{\partial t} \tilde{u}_i = -\frac{Z_i e}{m_i} \nabla \tilde{\phi} - \frac{T_i}{m_i n_{i0}} \nabla \tilde{n}_i. \]  

(II.1.5)

Taking the partial time derivative of eqn. II.1.4, the divergence of eqn. II.1.5 and combining them to eliminate the derivative of the ion velocity fluctuation, the equation for the density fluctuation of the i-th ion species \((1 \leq i \leq N)\) becomes

\[ \frac{\partial^2}{\partial t^2} \tilde{n}_i - \frac{T_i}{m_i} \nabla^2 \tilde{n}_i = \frac{Z_i n_{i0} e}{m_i} \nabla^2 \tilde{\phi}. \]  

(II.1.6)

Making \(A_1 = A_1 \cos(k_{0x}x - \omega_0 t)\) and Fourier-transforming equations II.1.2, II.1.3 and II.1.6, leads to a system of \((N + 2)\) equations—linear in the electron and ion species densities, and field potential fluctuations—,

\[ \hat{\phi}(k, \omega) - \frac{T_e}{en_{e0}} \hat{n}_e(k, \omega) = \frac{e}{2m_e c^2} A_1 \left[ \hat{A}(k - k_0, \omega - \omega_0) + \hat{A}(k + k_0, \omega + \omega_0) \right] \]

\[ - \frac{k^2}{4\pi e} \hat{\phi}(k, \omega) - \hat{n}_e(k, \omega) + \sum_{1 \leq i \leq N} Z_i \hat{n}_i(k, \omega) = 0 \]

(II.1.7)

\[ - \frac{Z_i n_{i0} e}{m_i} \frac{k^2}{(\omega^2 - k^2 u_{th,i}^2)} \hat{\phi}(k, \omega) + \hat{n}_i(k, \omega) = 0, \]

where the thermal speed for the i-th ion species is

\[ u_{th,i} = \sqrt{\frac{T_i}{m_i}}; \quad 1 \leq i \leq N. \]  

(II.1.8)

Eqn. II.1.7 can be written in matrix form
\[
\begin{bmatrix}
1 & -\frac{T_c}{\epsilon n_{c0}} & 0 & 0 & \ldots & 0 \\
-\frac{k^2}{4\pi e} & -1 & Z_1 & Z_2 & \ldots & Z_N \\
-\frac{Z_1n_{i0}e}{m_1} & \frac{k^2}{\omega^2 - k^2 u_{th,1}^2} & 0 & 1 & 0 & \ldots & 0 \\
-\frac{Z_2n_{i0}e}{m_2} & \frac{k^2}{\omega^2 - k^2 u_{th,2}^2} & 0 & 0 & 1 & \ldots & 0 \\
\vdots & \vdots & \vdots & \vdots & \ddots & \ddots & \vdots \\
-\frac{Z_Nn_{i0}e}{m_N} & \frac{k^2}{\omega^2 - k^2 u_{th,N}^2} & 0 & 0 & 0 & \ldots & 1 \\
\end{bmatrix}
\begin{bmatrix}
\phi \\
\hat{n}_c \\
\hat{n}_1 \\
\vdots \\
\hat{n}_N \\
\end{bmatrix}
= \begin{bmatrix}
0 \\
0 \\
0 \\
\vdots \\
0 \\
\end{bmatrix}.
\]

(II.1.9)

where
\[
\hat{A}^- = \hat{A}(k - k_0, \omega - \omega_0),
\]
\[
\hat{A}^+ = \hat{A}(k + k_0, \omega + \omega_0).
\]

When there is no illumination, the right hand side of eqn. II.1.9 is a null vector, and the determinant of the left hand side matrix provides a polynomial equation of order \(N\) in \(\omega^2\),

\[
1 - \frac{k^2}{1 + k^2 \lambda_{De}^2} \sum_{i=1}^{N} \frac{\omega^2}{\omega_{pi}^2 - k^2 u_{th,i}^2} = 0,
\]

(II.1.10)

where
\[
\lambda_{De} = \sqrt{\frac{T_c}{4\pi m_e e^2}} = \frac{u_{th,e}}{\omega_{pe}},
\]

(II.1.11)
is the electron Debye length, and
\[
\omega_{pi} = \sqrt{\frac{4\pi n_i Z_i^2 e^2}{m_i}},
\]

(II.1.12)
is the \(i\)-th (\(1 \leq i \leq N\)) ion plasma frequency.

This implicit polynomial in \(\omega^2\) is the dispersion relationship for the ion-acoustic modes present in a homogeneous, stationary, multispecies plasma. Solving this equation for \(\omega\), a set of \(N\) solutions of the form

\[
\omega_q = k_{ia} c_{sQ}(C_{s1}, C_{s2}, \ldots, C_{sN}; u_{th,1}, u_{th,2}, \ldots, u_{th,N}) \ (1 \leq q \leq N)
\]
can be found for any given value of \( k_{ia} \). These solutions—\( \omega_q = k_{ia} c_{sq} \) (1 \( \leq q \leq N \)—are the N ion-acoustic (low frequency) oscillation modes for a N-species plasma. A multispecies plasma has as many ion-acoustic modes as isotopic species are present. These modes are characterized by the value of the effective ion-acoustic phase speed for the \( q \)-th mode \( c_{sq} \) (1 \( \leq q \leq N \)). It depends only on the N values of the thermal velocities \( u_{th,i} \) (1 \( \leq i \leq N \)) of the N different ion species present in the plasma, and on the N values of their single-species ion-acoustic speeds—and therefore implicitly on the electron temperature \( T_e - C_{Si} \) (1 \( \leq i \leq N \)), defined as

\[
C_{si}^2 = \frac{\lambda_{De}^2 \omega_{pi}^2}{1 + k_{ia}^2 \lambda_{De}^2} + u_{th,i}^2 = \frac{Z_i T_e / m_i}{1 + k_{ia}^2 \lambda_{De}^2} + \frac{T_i}{m_i}.
\]  

(II.1.13)

For a single species plasma, the effective ion-acoustic phase speed of the only possible mode is

\[
c_s = C_{si} = \sqrt{\frac{\lambda_{De}^2 \omega_{pi}^2}{1 + k_{ia}^2 \lambda_{De}^2} + u_{th,i}^2}.
\]  

(II.1.14)

If one of the species has several ionization states present, only the effective ion-acoustic phase speed of the mode changes, but no additional modes appear due to charge effects. For example, a plasma composed by ion of a single species—but with M different states of ionization \( Z_j \) (1 \( \leq j \leq M \)) with ion number densities \( n_j \)—has M different values for the associated ion plasma frequencies \( \omega_{pj} \) (1 \( \leq j \leq M \)). But the value of the ion thermal speeds are independent of any of the M ionization states present in the plasma, so there is only a single value for the ion thermal speed \( u_{th} \) for a single isotopic species. The dispersion relationship II.1.10 then becomes,

\[
1 - \frac{k_{ia}^2 \lambda_{De}^2}{1 + k_{ia}^2 \lambda_{De}^2} \sum_{1 \leq j \leq M} \frac{\omega_{pj}^2 (n_{j0}, Z_j)}{\omega^2 - k_{ia}^2 u_{th}^2} = 0.
\]  

(II.1.15)

The denominator is independent of charge state

Therefore, the dispersion relationship II.1.10 has a single solution

\[
\omega^2 = k_{ia}^2 \left[ u_{th}^2 + \frac{\lambda_{De}^2}{1 + k_{ia}^2 \lambda_{De}^2} \sum_{1 \leq j \leq M} \omega_{pj}^2 \right] = k_{ia}^2 C_{SM}^2 \sum_{1 \leq j \leq M} \frac{n_{j0} Z_j^2}{n_{M0} Z_M^2}.
\]  

(II.1.16)
where $C_{sM}$ is now the ion-acoustic speed for a single species plasma with all its atoms fully ionized, with $n_{i0}$ the ion number density and $Z_M$ the charge state of the fully ionized atom. A similar result can be derived for a plasma containing more than one isotopic species, each isotopic species undergoing multiple states of ionization. The appearance of multimode ion-acoustic waves is strictly related to ion inertial effects.

The solution of the multi-isotopic linear system eqn. II.1.9. for the electron fluctuation $\hat{n}_e(k,\omega)$ when the laser light is present is

$$\hat{n}_e(k,\omega) = \frac{k^2}{8\pi m_c c^{2}} \left[ -1 + \sum_{l:\lambda_{l}=N} \frac{\omega_{pi}^2}{\omega^2 - k^2 \lambda_{l}^{2}} \right] \left[ \hat{A}(k-k_0,\omega-\omega_0) + \hat{A}(k+k_0,\omega+\omega_0) \right] A_l. (II.1.17)$$

To find the dispersion relationship for SBS in terms of the unperturbed parameters, all the frequency components of the electron density fluctuation $\hat{n}_e(k,\omega)$ and the scattered EM field vector potential fluctuation $\hat{A}(k,\omega)$ must be eliminated. As eqn. II.1.1 includes these fluctuations also, taking again $A_l = A_1 \cos(k_0 \cdot x - \omega_0 t)$, and Fourier-transforming eqn. II.1.1, we get

$$D(k,\omega) \hat{A}(k,\omega) = \frac{4\pi e^2}{m_e} \frac{A_1}{2} \left[ \hat{n}_e(k-k_0,\omega-\omega_0) + \hat{n}_e(k+k_0,\omega+\omega_0) \right], (II.1.18)$$

where

$$D(k,\omega) = \omega^2 - k^2 c^2 - \omega_{pe}^2. (II.1.19)$$

Evaluating the scattered EM wave vector potential at the down- and upshifted frequencies using eqn. II.1.18, gives

$$\hat{A}(k-k_0,\omega-\omega_0) = \frac{4\pi e^2}{m_e} \frac{A_1}{2} \frac{\hat{n}_e(k-2k_0,\omega-2\omega_0) + \hat{n}_e(k,\omega)}{D(k-k_0,\omega-\omega_0)}$$

$$\hat{A}(k+k_0,\omega+\omega_0) = \frac{4\pi e^2}{m_e} \frac{A_1}{2} \frac{\hat{n}_e(k,\omega) + \hat{n}_e(k+2k_0,\omega+2\omega_0)}{D(k+k_0,\omega+\omega_0)} \quad (II.1.20a,b)$$

Because the ion-acoustic mode is a low frequency mode ($\omega < \omega_0$) and non-resonant, the terms with $\hat{n}_e(k \pm 2k_0,\omega \pm 2\omega_0)$ can be neglected, as they do not satisfy the EM wave
dispersion relationship. Replacing the values of the down- and upshifted vector potentials, eqn. II.1.20a,b, in eqn. II.1.17 results in the dispersion relationship,

\[
1 + k^2 \lambda_{\text{De}}^2 \left[ 1 - \sum_{l \in \mathbb{Z}^N} \frac{\omega_{\text{pi}}^2}{\omega^2 - k^2 u_{\text{th},i}^2} \right] = \frac{k^2 v_{\text{os}}^2}{4} \left[ \frac{1}{D(k - k_0, \omega - \omega_0)} + \frac{1}{D(k + k_0, \omega + \omega_0)} \right],
\]

(II.1.21)

for the Stimulated Brillouin Scattering (SBS) of light in a multispecies, stationary homogeneous plasma. The electron quiver velocity is defined as

\[
v_{\text{os}} = \frac{e A_{\|}}{m_e c} = \frac{e E_{\parallel}}{m_e \omega_0}.
\]

(II.1.22)

In more practical units

\[
v_{\text{os}} = \frac{e}{m_e \omega_0} \sqrt{\frac{8 I_{\|}}{c}} = 26 \cdot 10^7 \sqrt{\frac{I_{\| [W/cm^2]}}{\lambda_{\text{De}} [cm]}} [\text{cm/s}].
\]

(II.1.23)

For Brillouin back- and sidescattering, where \( k \) is in the order of \( 2k_0 \), the upshifted light wave term can be dropped. Eqn. II.1.21 then becomes

\[
1 + k^2 \lambda_{\text{De}}^2 \left[ 1 - \sum_{l \in \mathbb{Z}^N} \frac{\omega_{\text{pi}}^2}{\omega^2 - k^2 u_{\text{th},i}^2} \right] = \frac{k^2 v_{\text{os}}^2}{4D(k - k_0, \omega - \omega_0)} \left[ 1 + \sum_{l \in \mathbb{Z}^N} \frac{\omega_{\text{pi}}^2}{\omega^2 - k^2 u_{\text{th},i}^2} \right]
\]

(II.1.24)

which, when multiplied by \( \prod_{l \in \mathbb{Z}^N} \left( \omega^2 - k^2 u_{\text{th},j}^2 \right) \) and \( D(k - k_0, \omega - \omega_0) \), becomes

\[
\text{LHS}(k, \omega)D(k - k_0, \omega - \omega_0) = \frac{k^2 v_{\text{os}}^2}{4} \frac{\text{RHS}(k, \omega)}{1 + k^2 \lambda_{\text{De}}^2}.
\]

(II.1.25)

LHS and RHS are

\[
\text{LHS}(k, \omega) = \left( 1 + k^2 \lambda_{\text{De}}^2 \right) \prod_{l \in \mathbb{Z}^N} \left( \omega^2 - k^2 u_{\text{th},j}^2 \right) - k^2 \lambda_{\text{De}}^2 \sum_{l \in \mathbb{Z}^N} \omega_{\text{pi}}^2 \prod_{j \neq i \in \mathbb{Z}^N} \left( \omega^2 - k^2 u_{\text{th},j}^2 \right),
\]

(II.1.26a)

and
\[
RHS(k, \omega) = - \prod_{l \leq j \leq N} \left( \omega^2 - k^2 u_{th,j}^2 \right) + \sum_{l \leq i \leq N} \omega_{pi}^2 \prod_{j \neq i} \left( \omega^2 - k^2 u_{th,j}^2 \right),
\]

(II.1.26b)

respectively. Equation II.1.25 is a polynomial of degree \(2N+2\) that must be solved numerically. For the single species case, the SBS dispersion relation is the quartic

\[
\left[ \omega^2 - k^2 u_i^2 - \frac{k^2 \lambda_{De}^2}{1 + k^2 \lambda_{De}^2} \omega_{pi}^2 \right] \mathcal{D}(k - k_0, \omega - \omega_0) = \frac{k^2 v_{rms}^2}{4} \left[ -\omega^2 + k^2 u_i^2 + \omega_{pi}^2 \right] \left( 1 + k^2 \lambda_{De}^2 \right). 
\]

(II.1.27)

There are two situations in which the solution to eqn. II.1.25 can be approximated. In the weak field limit, \(\omega_q = k_{iaq} c_{sq} \pm i \gamma_q\), with \(\gamma_q \ll k_{iaq} c_{sq}\), where \(q (1 \leq q \leq N)\) labels one of the \(N\) ion modes of the plasma. In these cases the solution reduce to one similar to Krueer's book.\(^9\) For a single species plasma, the resonant mode wavenumber \(k_{iaq}\) is\(^9\)

\[
k_{iaq} = 2k_0 \left( \cos \alpha - \frac{1}{\sqrt{1 - n_e/n_{\sigma}}} \frac{c_{sq}}{c} \right). 
\]

(II.1.28)

The growth rate \(\gamma_q\) associated with the ion-acoustic mode \(c_{sq}\), resonant at wavenumber \(k_{iaq}\) is

\[
\gamma_q^2 = \frac{k_{iaq}^2 v_{rms}^2}{16 \omega_0^3 k_{iaq}^3 c_{sq}^3} \sum_{l \leq i \leq N} \left( \frac{1 + k_{iaq}^2 \lambda_{De}^2}{c_{sq}^2 - u_{th,i}^2} \right)^2 \left[ \frac{1}{1 + k_{iaq}^2 \lambda_{De}^2} \sum_{l \leq i \leq N} \frac{1}{c_{sq}^2 - u_{th,i}^2} \right]. \]

(II.1.29)

As there are \(N\) solutions \(\{(k_{iaq}, c_{sq}, \gamma_q), (1 \leq q \leq N)\}\) to the SBS dispersion equation II.1.25, only the solution with the highest growth rate,

\[
\gamma = \text{Max}\{\gamma_q, (1 \leq q \leq N)\}, 
\]

(II.1.30)

must be considered. For a single species, homogeneous, stationary plasma and with \(k_{iaq} \lambda_{De} \ll 1\), eqn. II.1.31 reduces to the usual expression\(^10\)

\[
\gamma^2 = \frac{k_0^2 v_{rms}^2 \omega_{pi}^2}{8 \omega_0^3 c_i^3} \cos \alpha. 
\]

(II.1.31)

The resonant mode label \(q (1 \leq q \leq N)\) will be dropped when referring to the elements of the fastest growing solution of the SBS dispersion equation II.1.21, which will be expressed as \((k_{ia}, c_i, \gamma)\).
Fig. II.1- A wavenumber vector diagram of SBS.

\( \alpha \) is the angle of the ion-acoustic wavenumber vector \( \mathbf{k}_{ia} \) and \( \beta \) is the Brillouin scattering angle of the electromagnetic wavenumber vector \( \mathbf{k}_s \), both taken with respect to the incident laser wavenumber vector direction \( \mathbf{k}_0 \).

The expression II.1.29 for the SBS growth rate \( \gamma \) predicts that for identical plasma density and temperature conditions, illuminated by identical incident laser intensities, the ion-acoustic wave associated with the backscattered light \((\alpha = 0, \beta = \pi)\) grows faster than the ion-acoustic wave associated with the sidescattered light \((\alpha > 0, \beta < \pi)\)

\[ \gamma \text{ (back)} \geq \gamma \text{ (side)}. \]

In the **strong field limit**, \( |\omega_i| \gg |k_{ia}c_s| > k_{ia}u_{th,i} \) \((1 \leq i \leq N)\). Then eqn. II.1.25 reduces to the polynomial

\[ \omega^2 D(k - k_0, \omega - \omega_0) = \frac{k_{ia}^2 v_{ox}^2}{4} \left( -\omega^2 + \sum_{i=1}^{N} \omega_{pi}^2 \right). \tag{II.1.32} \]

As

\[ D(k - k_0, \omega - \omega_0) = \omega^2 - 2\omega_0 \omega - k_{ia}^2 c^2 + 2k_0 \cdot k_{ia} c^2 + \omega_0^2 - k_{0}^2 c^2 + \omega_{pe}^2 \]

\[ = -2\omega_0 \omega - k_{ia}^2 c^2 + 2k_{ia} k_0 c^2 \cos \alpha. \]

with \( \omega \ll \omega_0 \), then taking \( k_{ia} = 2k_0 \cos \alpha \), eqn. II.1.33 becomes the cubic

\[ \omega^3 = \frac{k_{ia}^2 v_{ox}^2}{2\omega_0} \left( \omega^2 - \sum_{i=1}^{N} \omega_{pi}^2 \right) \cos^2 \alpha. \tag{II.1.33} \]
CHAPTER II - THEORY

This equation has a solution

\[ \omega = \frac{k_0^2 v_{in}^2}{2\omega_0} \cos^2 \alpha \left\{ \frac{1}{3} + \frac{1 + i\sqrt{3}}{6} \left[ \Xi - 1 + \sqrt{(\Xi - 1)^2 - 1} \right] \right\}^{\frac{1}{3}} + \frac{1 - i\sqrt{3}}{6} \left[ \Xi - 1 - \sqrt{(\Xi - 1)^2 - 1} \right] \right\}^{\frac{1}{3}}. \] (II.1.34)

where

\[ \Xi = 54 \frac{\omega_0^2 \cos^{-1} \alpha}{k_0 v_{in}^2} \sum_{l=1}^{N} \omega_p^2. \]

As usually \( \Xi >> 1 \), the solution for the strong field limit is

\[ \omega_{ia} = \left( \frac{k_0^2 v_{in}^2}{2\omega_0} \cos^2 \alpha \sum_{l=1}^{N} \omega_p^2 \right)^{\frac{1}{3}} \left[ 1 + \frac{i\sqrt{3}}{6} \right], \] (II.1.35)

with the growth rate \( \gamma = \text{Im}[\omega_{ia}] \).

Even for a multispecies plasma, in the strong field limit there is only a single SBS mode. In the strong field limit—similarly to the weak field case—for identical plasma density and temperature conditions, illuminated by identical incident laser intensities, the ion-acoustic wave associated with the backscattered light (\( \alpha = 0, \beta = \pi \)) grows faster than the ion-acoustic wave associated with the sidescattered light (\( \alpha > 0, \beta < \pi \))

\[ \gamma \ (\text{back}) \geq \gamma \ (\text{side}). \]

Figure II.2a-c shows comparative plots of the full quartic and strong field limit solutions of equation II.1.25, \( \text{Re}[\omega_{ia}] \) and \( \text{Im}[\omega_{ia}] \), for Brillouin backscattering in fully ionized Silicon (\( A = 28, Z = 14 \)), at some representative electron number densities.

B- The Brillouin Scattered Light Spectrum

The results in the previous section have been limited to describe the ion-acoustic wave \((k_{ia}, \omega_{ia})\) interacting through the Brillouin scattering with an incident electromagnetic wave \((k_0, \omega_0)\). The parameters of the ion-acoustic wave can not be measured directly by an observer external to the interacting region inside the plasma. To probe the Brillouin interaction, only the spectral characteristics of the scattered electromagnetic wave \((k_s, \omega_s)\) are available. The most readily observable of these spectral parameters is its wavelength \(\lambda_s\). Using the solution to the polynomial equation II.1.24 or the strong-field limit solution II.1.35 and the Manley-Rowe relation I.1 the scattered wavelength is given by
Fig. II.2a,b- Quartic and strong-field solutions to the Brillouin dispersion relation.
For this plasma, $T_e=10$ eV, $T_e=100$ eV, and $\lambda_L=1054$ nm.
Fig. II.2c - Quartic and strong-field solutions to the Brillouin dispersion relation.
For this plasma, \( T_e = 10 \text{ eV}, \ T_i = 100 \text{ eV}, \) and \( \lambda_L = 1054 \text{ nm}. \)

\[
\omega_s = \omega_0 - \text{Re}[\omega_{ia}]. \tag{II.1.37}
\]

The Brillouin scattered light seen by an observer outside of the plasma wavelength is

\[
\lambda_s = \frac{2\pi c}{\omega_s} = \frac{2\pi c}{\omega_0 - \text{Re}[\omega_{ia}]} = \frac{\lambda_l}{1 - \lambda_l \cdot \text{Re}[\omega_{ia}]/2\pi c}. \tag{II.1.38}
\]

The absolute Brillouin shift in a homogeneous, stationary plasma is

\[
\frac{\Delta \lambda_{\text{SBS}}}{\lambda_l} = \frac{\lambda_s - \lambda_l}{\lambda_l} = \frac{1}{2\pi c/\lambda_l \cdot \text{Re}[\omega_{ia}]} - 1 = \frac{\lambda_l \cdot \text{Re}[\omega_{ia}]}{2\pi c}. \tag{II.1.39}
\]

as usually \( \lambda_l \cdot \text{Re}[\omega_{ia}]/2\pi c \ll 1. \)

The Brillouin scattered light wavelength \( \lambda_s \) is always redshifted with respect to the incident laser wavelength \( \lambda_l. \) For identical laser beam intensity, and plasma density and temperature conditions, the backscattered light (\( \alpha = 0, \ \beta = \pi \)) is redshifted more than the sidescattered light (\( \alpha > 0, \ \beta < \pi \)), as the backscattered SBS mode frequency \( \text{Re}[\omega_{ia}] \) is
approximately \( \cos^{-2/3} \alpha \) higher than the sidescattered mode SBS frequency. Hence, in a stationary plasma is

\[
\lambda_\text{s (back)} \leq \lambda_\text{s (side)} \leq \lambda_\text{l},
\]

### C- The SBS Threshold with Collisional and Landau Damping

The damping of the unstable scattered EM and ion-acoustic waves introduces a threshold intensity for the onset of the stimulated Brillouin instability. Incorporating the damping terms \( v_{IA} \partial n_A/\partial t \) in the ion density fluctuation eqn. II.1.6 and \( v_{FM} \partial A/\partial t \) in the EM wave coupling eqn. II.1.1—with effective energy damping rates \( v_{IA} \) for the ion-acoustic wave and \( v_{FM} \) for the electromagnetic wave—, the SBS dispersion equation II.1.20 has the same dependencies with the substitutions

\[
\omega^2 - u_{\text{th,i}}^2 \rightarrow \omega(\omega + i v_{IA}) - u_{\text{th,i}}^2 \quad \text{(II.1.40)}
\]

and,

\[
D(k, \omega) = \omega^2 - k^2 c^2 + \omega_{pe}^2 \rightarrow \omega(\omega + i v_{FM}) - k^2 c^2 + \omega_{pe}^2.
\quad \text{(II.1.41)}
\]

Hence, the dispersion relationship for downshifted SBS, eqn. II.1.24, becomes

\[
1 + k^2 \lambda_{1C}^2 \left[ 1 - \sum_{l \leq N} \frac{\omega_{\text{pi}}^2}{\omega(\omega + i v_{IA}) - k^2 u_{\text{th,i}}^2} \right] = \frac{k^2 v_{\text{in}}^2/4}{(\omega - \omega_0)(\omega - \omega_0 - iv_{FM}) - (k - k_\alpha)^2 c^2 - \omega_{pe}^2}
\quad \text{(II.1.42)}
\]

The new damped solutions \( \gamma^* \) of eqn. II.1.42 are found using a procedure similar to that used to find the undamped solutions \( \gamma \) of eqn. II.1.21. These solutions \( \gamma^* \) are related to the undamped solutions \( \gamma \) of eqn. II.22 by

\[
\left( \gamma^* + v_{FM}/2 \right) \left( \gamma^* + v_{IA}/2 \right) = \gamma^2.
\quad \text{(II.1.43)}
\]

The threshold condition for the growth rate with damping of both the ion-acoustic and electromagnetic waves is

\[
\gamma \geq \sqrt{v_{FM} v_{IA}/2}.
\quad \text{(II.1.44)}
\]
Typically, the electromagnetic wave energy decays by collisional damping and the ion-acoustic wave by Landau damping, raising the intensity threshold for the appearance of the Brillouin instability.\textsuperscript{11}

Collisional damping (also commonly referred to as inverse-bremsstrahlung) dissipates the energy of the electromagnetic field. Collisions of electrons oscillating in the field with ions thermally spread part of the coherent kinetic energy acquired by the electrons interacting with the electromagnetic field. This damping reduces the strength of the ponderomotive force which creates the plasma density fluctuations through depletion of the scattered EM wave (depletion of the incident EM wave is negligible) and attenuates the backscattered light intensity. The electron-ion collision frequency for a multispecies plasma (with a Maxwellian temperature distribution for all species) is,\textsuperscript{12}

\[
v_{ci} = \frac{1}{3(2\pi)^{3/2}} \frac{\omega_{pe}^2}{\omega_0^2} \frac{\omega_{pe}^4}{n_e u_{th,e}^3} \sum_{i=1}^N \frac{n_i}{n_e} Z_i^2 \ln \Lambda_{ci} . \tag{II.1.45}
\]

where \( \Lambda_{ci} \) is the ratio of the maximum \( (\lambda_{ci}) \) to the minimum \( (Z_i e^2/m_e v_{th,e}^2) \) impact distances between electrons and ions; \( \ln \Lambda_{ci} \) is also known as the collisional Coulomb logarithm. It can be shown that\textsuperscript{13}

\[
\Lambda_{ci} = 12\pi n_e \lambda_{ci}^3/Z_i . \tag{II.1.46}
\]

Landau damping is a collisionless mechanism in which the ion-acoustic wave dissipates its energy by accelerating and decelerating particles of the plasma in resonance with the wave, moving at \( \gamma_{\text{particle}} = \omega_{ia}/k_{ia} \).\textsuperscript{14} This is a kinetic effect, and the hydrodynamic form of the dispersion relationship II.1.10 is not sufficient to describe the behavior of the ion-acoustic wave completely.

If it is assumed that \( \text{Re}[\omega_{ia}] >> \text{Im}[\omega_{ia}] \), an elegant calculation of the Landau damping of an ion-acoustic wave starts by writing its dispersion relationship in terms of the plasma dielectric function,\textsuperscript{15}

\[
\varepsilon(k_{ia}, \omega_{ia}) = \text{Re}[\varepsilon(k_{ia}, \omega_{ia})] + i\text{Im}[\varepsilon(k_{ia}, \omega_{ia})] \approx 0 . \tag{II.1.47}
\]

Taylor expanding about \( \omega_{ia} = \text{Re}[\omega_{ia}] \) yields
\[ \text{Re} \left[ \varepsilon(k_{ia}, \text{Re} \omega_{ia}) \right] + i \text{Im} \left[ \varepsilon(k_{ia}, \text{Re} \omega_{ia}) \right] + i \text{Im} \left[ \omega_{ia} \right] \frac{\partial \text{Re} \left[ \varepsilon(k_{ia}, \omega_{ia}) \right]}{\partial \omega_{ia}} \bigg|_{\omega_{ia} = \text{Re} \omega_{ia}} = 0. \]  

(II.1.48)

As both factors in the product

\[ \text{Im} \left[ \omega_{ia} \right] \frac{\partial \text{Im} \left[ \varepsilon(k_{ia}, \omega_{ia}) \right]}{\partial \omega_{ia}} \bigg|_{\omega_{ia} = \text{Re} \omega_{ia}} \]

are small, this term is ignored. Equating both the real and imaginary parts of II.1.48 to zero, then

\[ \text{Re} \left[ \varepsilon(k_{ia}, \text{Re} \omega_{ia}) \right] = 0, \]  

(II.1.49)

and

\[ \text{Im} \left[ \omega_{ia} \right] = -\frac{\text{Im} \left[ \varepsilon(k_{ia}, \omega_{ia}) \right]}{\frac{\partial}{\partial \omega_{ia}} \text{Re} \left[ \varepsilon(k_{ia}, \omega_{ia}) \right]} \bigg|_{\omega_{ia} = \text{Re} \omega_{ia}}. \]  

(II.1.50)

The kinetic dispersion relationship II.1.49 of the ion-acoustic wave, is equivalent to the hydrodynamic equation II.1.10, developed in section II.1A. Finding the solutions of eqn. II.1.49 gives the multiple ion-acoustic modes of the plasma, but in this formulation a more detailed (but computationally more taxing) kinetic description is achieved.

The plasma dielectric function can be written in its kinetic form (indexes \( i > 0 \) label the \( N \) ion species, \( i = 0 \) labels the electrons) as\(^{16}\)

\[ \varepsilon(k_{ia}, \omega_{ia}) = 1 + \sum_{0 < i < N} \frac{1}{k_{ia}^2 \lambda_{ia}^2} \left[ 1 + \zeta_i Z(\zeta_i) \right]. \]  

(II.1.51)

\( Z(\zeta) \) is the plasma dispersion function,\(^{17}\) defined (for all \( \zeta \) in the complex plane) as

\[ Z(\zeta) = 2i e^{-\frac{\zeta^2}{2}} \int_{-\infty}^{\zeta} e^{-t^2} dt. \]  

(II.1.52)

where \( \zeta_{si} \) is the ratio of the ion-acoustic phase speed to the thermal speed,

\[ \zeta_{si} = \frac{\omega_{ia}}{\sqrt{2k_{ia} u_{th,i}}} = \frac{c_s}{\sqrt{2} u_{th,i}}. \]  

(II.1.53)
Taking the asymptotic expansion (\( \zeta_i \approx 1 \), for the ions) and Taylor series (\( \zeta_0 \ll 1 \), for the electrons) in eqn. II.1.51, the real and imaginary parts of the plasma dielectric function are

\[
\mathcal{R}[\varepsilon(k_{ia},\omega_{ia})] = 1 + \frac{\omega_{pe}^2}{k_{ia}^2 u_{th,c}^2} - \sum_{l=1}^{\infty} \frac{2\omega_{pi}^2}{k_{ia}^2 u_{th,i}^2 \zeta_l} = 1 + \frac{\omega_{pe}^2}{k_{ia}^2 u_{th,c}^2} - \sum_{l=1}^{\infty} \frac{\omega_{pi}^2}{\omega_{ia}^2} \tag{II.1.54}
\]

\[
\mathcal{I}[\varepsilon(k_{ia},\omega_{ia})] = \sqrt{\pi} \sum_{l=1}^{\infty} \frac{\omega_{pi}^2}{k_{ia}^2 u_{th,i}^2} e^{-\zeta_i^2} \tag{II.1.55}
\]

Applying eqn. II.1.50, the Landau damping frequency is given by

\[
u_{1an} = \mathcal{I}[\omega_{ia}] = \sqrt{\pi} \frac{\sum_{l=1}^{\infty} \omega_{pi}^2 \zeta_i^2 e^{-\zeta_i^2}}{\sum_{l=1}^{\infty} \omega_{pi}^2} \bigg|_{\omega_{ia} = \mathcal{R}[\omega_{ia}]} \tag{II.1.56}
\]

or more explicitly, for a Maxwellian plasma,

\[
u_{1an} = \frac{\sqrt{\pi}}{8} k_{ia} c_s^4 \frac{\sum_{l=1}^{\infty} n_{l0} Z_i^2 e^{-c_s^2/2u_{th,i}^2}}{\sum_{l=1}^{\infty} \frac{n_{l0} Z_i^2}{m_l}} \tag{II.1.57}
\]

Making the identification \( v_{1M} = v_{1c} \) and \( v_{1A} = v_{1an} \) in equation II.1.43, the presence of dissipative processes, collisional and Landau damping, leads to a SBS convective threshold condition

\[
\gamma \geq \sqrt{\frac{v_{1c}}{v_{1an}}} / 2. \tag{II.1.58}
\]

The minimum laser intensity needed to overcome the effects of collisional and Landau damping (for Brillouin backscattering in a single species homogeneous Maxwellian, non-drifting plasma) is the dissipative (convective) SBS threshold

\[
I_{SBS}^{\text{diss}} > \frac{1}{12\pi} m_c \left( \frac{m_i}{m_c} \right)^2 \omega_0^3 \frac{n_c}{n_\sigma} \frac{c}{u_{th,c}^3} \left[ e^{-c_s^2/2u_{th,i}^2} + \frac{\omega_{pi}^2}{\omega_{pe}^2} \left( \frac{u_{th,i}^3}{u_{th,c}^3} \right) e^{-c_s^2/2u_{th,i}^2} \right] \ln \Lambda \tag{II.1.59}
\]

In more practical units, the dissipative (convective) SBS threshold \( I_{SBS}^{\text{diss}} \) is
\[
I_{\text{diss}}^{\text{SBS}} \left[ \frac{\text{W}}{\text{cm}^2} \right] > \frac{5.4 \cdot 10^{11}}{A_i \lambda_{\text{i,\{\mu m}\}}^3 T_i^{\text{f/eV}^2} n_{\sigma}} \left( \frac{Z_i T_e/T_i + 1}{T_c/T_i} \right)^{3/2} \ln \left[ 0.46 \frac{\lambda_{\text{i,\{\mu m}\}}}{Z_i} \frac{T_i}{\sqrt{n_c/n_{\sigma}}} \right] \times \left[ \exp \left( -\frac{1}{1836 A_i} \right) \left( \frac{T_i}{Z_i T_e/T_i + 1} \right) \right] + Z_i \left( \frac{T_i}{Z_i T_e/T_i + 1} + 1 \right) \right),
\]

(II.1.60)

Figure II.3 shows the dissipative SBS threshold intensity as a function of the electron temperature \(T_e\), in different temperature ratios \(T_e/T_i\), for a homogeneous stationary fully-ionized Silicon \((Z_i = 14, \ A_i = 28)\) plasma, illuminated with normally incident infrared light \((\lambda_i = 1.054 \ \mu\text{m})\), and for \(n_c = 0.1 \ n_{\text{cr}}\). For a fully ionized silicon plasma, with an electron number density \(n_c = 0.1 \ n_{\text{cr}}\), electron temperature \(T_e = 1 \ \text{keV}\) and ion temperature \(T_i = 0.1 \ \text{keV}\), the dissipative thresholds for infrared \((1054 \ \text{nm})\) and ultraviolet \((351 \ \text{nm})\) are

\[
I_{\text{diss}}^{\text{SBS}} > \begin{cases} 
2.5 \cdot 10^{12} \ \text{W/cm}^2, \quad \lambda_{\text{i,}} = 1054 \ \text{nm} \\
6.7 \cdot 10^{13} \ \text{W/cm}^2, \quad \lambda_{\text{i,}} = 351 \ \text{nm}.
\end{cases}
\]

**Fig. II.3 - Dissipative SBS threshold intensity for a homogeneous stationary fully-ionized Si plasma \((n_c/n_e = 0.1)\) for normally incident infrared light \((\lambda_L = 1.054 \ \mu\text{m})\).**

For temperature ratios \(T_e/T_i > 5\), the threshold plots (long dashed lines) are approximately the same.
In this experiment—using Silicon, Gold and Parylene-N—the expected values of the dissipative SBS threshold for $\lambda_i = 1.054 \, \mu m$ are well below the available intensities ($10^{14} \leq I_i \leq 5 \cdot 10^{16} \, W/cm^2$) for the TTT-CPA laser (see § III.2).

**D- Non-isothermal effects for high-Z plasmas.**

Ordinary theoretical analysis of SBS in laser produced plasmas assumes that the plasma is isothermal and Maxwellian. This assumption is justified by the fact that the ion-acoustic wavelength is very small (about half of the laser wavelength), so that, using classical (Spitzer-Harm\textsuperscript{18}) thermal transport theory, the time for thermal equilibration across this distance is much smaller than the ion-acoustic wave period. As seen in section II.1A, the driving term for the instability in this model is the ponderomotive force.

Recent numerical studies of thermal transport in laser-fusion plasmas using the Fokker-Planck (FP) equation have shown that classical transport theory is inadequate to treat phenomena occurring over short distances, even if the local temperature scalelength $T_e/V_T e$ is much longer than the electron mean free path.\textsuperscript{19,20,21} The thermal conductivity can be greatly reduced for temperature variations with wavelengths shorter than the mean free path of those electrons which, classically, carry the bulk of the heat flow. Here, electrons faster than those with thermal energies (with velocities $> 3.7 \, v_{th,e}$), rapidly distribute themselves uniformly and decouple from the spatial variations in energy density that dominate the slower electrons. The slower electrons contain most of the thermal energy. Thus, these local variations persist longer than predicted by the classical theory. This effect is especially significant if there is a source which preferentially heats the slower (thermal) electrons.

Collisions (inverse-bremsstrahlung) are the principal heating mechanism in laser produced plasmas. Collisional heating raises the temperature and pressure of the plasma in regions of high electromagnetic field intensity. The associated ponderomotive force can expel plasma out of this region. With classical thermal conductivity, the electron thermal energy spreads uniformly through the bulk of the plasma, thus making any local temperature variation negligible. With a ‘non-local’ transport theory these variations are significant and in high-Z (as gold) plasmas they can be the dominant driving force for the SBS instability.

Short and Epperlein\textsuperscript{22} analyzed the non-isothermal SBS instability for an stationary homogeneous equilibrium plasma, referred from here on as the non-local SE model. The electric field of the incident laser light is represented as
\[ E_L = E_0(t)e^{i(k_0x - \omega_0 t)} + E_0^\ast(t)e^{-i(k_0x - \omega_0 t)}, \]

and the electric field of the scattered light is
\[ E_S = \tilde{E}_x(t)e^{i(k_xx - \omega_x t)} + \tilde{E}_x^\ast(t)e^{-i(k_xx - \omega_x t)}. \]

The plasma has a uniform electron number density and temperature, which are perturbed by slowly varying fluctuations,
\[ n_c = n_{c0} + \tilde{n}_c(t)e^{i(k_{ia}x - \omega_{ia} t)} + \tilde{n}_c^\ast(t)e^{-i(k_{ia}x - \omega_{ia} t)}, \]
\[ T_c = T_{c0} + \tilde{T}_c(t)e^{i(k_{ia}x - \omega_{ia} t)} + \tilde{T}_c^\ast(t)e^{-i(k_{ia}x - \omega_{ia} t)}. \]

Backscattering is considered for simplicity and the Manley-Rowe matching conditions are assumed to be satisfied. As in the isothermal case, there is an equation for the coupling of the electromagnetic field to the electron fluid,
\[ 2i\omega_x \frac{\partial}{\partial t} \tilde{E}_x^\ast(t) = -\frac{\omega_x}{\omega_0} \frac{\omega_p^2}{\omega_0} E_0 \frac{\tilde{n}_c(t)}{n_{c0}}, \quad \text{(II.1.61)} \]

and an equation for the electron density fluctuation,
\[ \left[ \frac{\omega_{ia}^2 - k_{ia}^2 c_s^2 + 2i\omega_{ia}}{n_{c0}} \frac{\partial^2}{\partial t^2} - \frac{\partial^2}{\partial t^2} \right] \tilde{n}_c(t) = \frac{k_{ia}^2 c_s^2}{n_{c0}} \frac{\tilde{T}_c(t)}{T_{c0}} + \frac{Ze^2}{m_e m_i \omega_0 \omega_x} E_0 \tilde{E}_x^\ast(t). \quad \text{(II.1.62)} \]

Balancing collisional heating with thermal diffusion, and taking into account the temperature and density dependence of the collisional absorption coefficient, the expression for the temperature variation is
\[ \left[ \frac{3}{2} \left( \frac{\partial}{\partial t} - i\omega_{ia} \right) + \kappa_0 \frac{k_{ia}^2}{n_{c0}} \frac{1}{T_{c0}} \right] \tilde{T}_c(t) = 2 \frac{\kappa_{0e} l_1}{n_{c0} T_{c0}} \frac{\tilde{n}_c(t)}{n_{c0}} + \frac{ce_0^2 \kappa_{0e} l_1}{2 \pi n_{c0} T_{c0}} E_0 \tilde{E}_x^\ast(t). \quad \text{(II.1.63)} \]

The collisional and thermal diffusion coefficients \( \kappa_0^{\alpha} \), \( \kappa_0^{th} \) can be written in terms of the Schmitt\( ^{2,3} \) dimensionless parameters \( \gamma_{11}, \gamma_{12} \) and the classical Spitzer-Harm conductivity \( \kappa_0^{SH} \),
\[ \kappa_0^{\alpha} = \frac{3}{2} \frac{\omega_{ia} n_{c0} T_{c0}}{l_1} \frac{\gamma_{12}}{\gamma_{11}}, \quad \text{(II.1.64)} \]
\[ \kappa_0^{th} = \frac{3 \omega_{ia} n_{e0}}{2 k_{ia}^2 \gamma_{T11}} \kappa_0^{SH} \]  

(II.1.65)

\[ \gamma_{T1} = 6.75 \frac{\ln \Lambda_{e1}}{\varepsilon_0 \mu_0^2 \gamma_{T11}^2} \Phi(Z) \left( \frac{Z}{A} \right)^2 \frac{n_{e0}}{n_e}, \]  

(II.1.66)

and \( \gamma_{T2} \) is the ratio of the collisional heating rate to the thermal cooling rate across \( k_{ia} \):

\[ \gamma_{T2} = 2.24 \cdot 10^{-8} \frac{I_{L} \left[ \text{W/cm}^2 \right]}{e_0 \mu_0^2 \gamma_{T11}^2 \Phi(Z) \left( \frac{n_{e0}}{n_e} \ln \Lambda_{e1} \right)^2}. \]  

(II.1.67)

The ratio of the thermal conductivity for collisional heating, \( \kappa_0^{th} \), to the classical Spitzer-Harm conductivity, \( \kappa_0^{SH} \), is found numerically to be well approximated by

\[ \frac{\kappa_0^{th}}{\kappa_0^{SH}} = \frac{1}{1 + \left[ 84 \pi k_{ia} n_{e0} \lambda_{1e}^4 \sqrt{4.2Z / \Phi(Z) \ln \Lambda_{e1}} \right]^{1/44}}. \]  

(II.1.68)

Ion motion and time-varying heating have no significant effect on the value of this ratio, with

\[ \Phi(Z) = \frac{0.24 + Z}{1 + 0.24Z}. \]  

(II.1.69)

For a time dependence of the perturbed quantities of the form \( e^{i \Omega t} \), the equations II.1.61-63 become a set of simultaneous algebraic equations that can be combined in the quartic polynomial in \( \Omega \):

\[ \Omega \left[ (\omega_{ia} + \Omega)^2 - k_{ia}^2 c_s^2 \right] \left[ \frac{3}{2} i (\omega_{ia} + \Omega) - \frac{k_{ia}^2}{n_{e0}} - \frac{3 \kappa_0^{th} I_{L1}}{2 n_{e0} T_{e0}} \right] + \]

\[ \frac{4 \pi I_{11}}{c} \frac{\omega_{pe}^2}{\omega_0^2 \omega_{is}} k_{ia}^2 \frac{Z e^2}{m_e m_i} \left[ \frac{3}{2} i (k_{ia} c_s + \Omega) - \frac{k_{ia}^2}{n_{e0}} - \frac{3 \kappa_0^{th} I_{L1}}{2 n_{e0} T_{e0}} \right] + \]

\[ 2 k_{ia}^2 c_s^2 \frac{\kappa_0^{th} I_{L1}}{n_{e0} T_{e0}} \left[ \Omega - \frac{\omega_{pe}^2}{\omega_0} \frac{\nu^2}{\omega_0} \right] = 0. \]  

(II.1.70)
The complex value of the ion-acoustic frequency $\gamma = \text{Im}[\omega_{ia}]$ is given by equations II.1.10 and II.1.34. In all cases the growth rates are maximized for values of $k_{ia}$ near the resonance of the scattered EM wave

$$k_{ia} \approx 2k_0 \left[ 1 - \frac{1}{\sqrt{1 - n_e/n_0}} \frac{c_s}{c} \right].$$  \hspace{1cm} (II.1.71)

The maximum (positive) value of the imaginary parts of the four complex roots to eqn. II.1.63 gives the growth rate $\Gamma$ for the non-local SE model. Figure II.4 shows the square of the ratio of the non-local SE growth rate $\Gamma$ to the isothermal growth rate $\gamma$ for Gold ($Z = 79$, $A = 193$) and Silicon ($Z = 14$, $A = 28$), as a function of the electron temperature. This ratio represents the relative magnitude of the non-local thermal transport contribution to that of the ponderomotive force, and it is the same as the ratio of the laser intensity thresholds,

$$\frac{\Gamma^2}{\gamma^2} = \left( \frac{\Gamma_{\text{th,LS}}}{\Gamma_{\text{is,th}}} \right)^{-1} = 1 + \frac{U_{\text{th,c}}^2}{V_{\text{os}}^2} \gamma \Gamma \left[ \frac{k_{\text{th}}}{k_{\text{sh}}} \right]^{-1}. \hspace{1cm} (II.1.72)$$

For hot (> 500 eV) low-Z plasmas it is found that non-local thermal transport has a negligible effect, and the ponderomotive force remains the principal driving term for SBS. However, for high-Z plasmas collisional heating becomes significant, with the second term of eqn. II.1.72 providing the main contribution to the growth rate.

The non-isothermal term becomes more important for a larger ionization state $Z$ and electron number density $n_{e0}$, and diminishes with increasing electron temperature $T_{e0}$ and laser wavelength $\lambda_L$. Even for low-Z plasmas such as silicon, the contribution of the non-isothermal terms cannot be ignored for electron temperatures below the 500 eV range, while its effect is negligible for higher electron temperatures. A SBS threshold reduction up to two orders of magnitude can be expected for low temperature (< 100 eV) plasmas, for any value of the ionization state $Z$ of the plasma ions (see §IV.2A and §V.2 for the measured SBS thresholds).
Fig. II.4 - Ratio of the square of the SBS growth rates of the non-local transport model to the isothermal model for homogeneous stationary fully-ionized Au and Si plasmas.

The ratio of the square of the non-local SE model SBS growth rate to the isothermal SBS model growth rate $(\Gamma / \gamma)^2$ is the same as the ratio of the SE model laser intensity threshold to the isothermal model laser intensity threshold $I_{th}^{\text{SE}} / I_{th}^{\text{iso}}$. 

**Gold**

$A=197$ $Z=79$

Electron temperature, $T_{e0}$ [eV]

**Silicon**

$A=28$ $Z=14$

Electron temperature, $T_{e0}$ [eV]
II.2 - SBS IN A DRIFTING INHOMOGENEOUS PLASMA

A- The Doppler and Index of Refraction Shifting of the Backscattered Light

SBS is not the only source of frequency shifting in the observed backscattered light. As the plasma expands isothermally, the interaction between the incident laser light and the plasma occurs in the reference frame fixed within the region of the plasma that supports the resonance. This region may expand with a finite speed, which may be highly supersonic (faster than the local ion-acoustic speed), with respect to the static laboratory frame.

For an external observer, if the plasma is flowing towards the observer, the measured backscattered light has a large Doppler blueshift superimposed in the usual Brillouin scattering red shifting. The wavelength of the scattered light now has the expression:

\[
\lambda_s = \lambda_i \left[ 1 - \frac{u_{ex}}{c} \right] + \Delta \lambda_{SBS},
\]  

(II.2.1)

where \( u_{ex} \) is the plasma expansion speed (along the direction of the incident beam) for the SBS resonance region at position \( x_{res} \) (see the next section) where most of the emission occurs.

Also a spectral shift results from the change in the plasma's refraction index during the period between the arrival and departure of the interaction laser pulse into the SBS resonance region. The refraction index changes due to the variation of the electron number density \( n_e(x,t) \) along the plasma profile with its expansion meanwhile the interaction pulse propagates through it. A simple expression for this refraction index shift can be found,

\[
\frac{\Delta \lambda}{\lambda_i} = -\frac{1}{c} \frac{d}{dt} \left[ \int_{x > x_{res}} \sqrt{1 - \frac{n_e(x)}{n_{i\tau}}} \, dx \right],
\]  

(II.2.2)

where the integral is taken outwards from the position of SBS emission \( x_{res} \) along the plasma profile with its expansion meanwhile the interaction pulse propagates through it. Depending on the temporal evolution of the plasma density profile, the refraction index shift can be to the red or to the blue side of the spectrum. For both the planar and spherical isothermal models (see § III.1.B) for the expansion of a plasma the refraction index always blueshifts the incident light.
Using eqn. II.1.37, the Brillouin scattering spectral shift of the scattered EM wave (at angle $\beta$ with respect to the incident beam direction) is

$$\Delta \lambda_{\text{SBS}} = \lambda_{\text{i}} \frac{\lambda_{\text{i}} \text{Re}[\omega_{\text{ia}}]}{2 \pi c}. \quad (\text{II.2.3})$$

Therefore the total observed spectral shift will be

$$\Delta \lambda = \lambda_{\text{i}} \left\{ \frac{\lambda_{\text{i}} \text{Re}[\omega_{\text{ia}}]}{2 \pi c} - \frac{u_{\text{cr}}}{c} \frac{1}{c} \frac{\partial}{\partial t} \left[ \int_{x>x_{\text{res}}} \sqrt{1 - \frac{n_{\text{e}}(x)}{n_{\text{cr}}}} \, dx \right] \right\}. \quad (\text{II.2.4})$$

The first term represents the usual SBS redshift, the second term is the Doppler blueshift produced by the motion of the expanding plasma towards the observer, and the third term is the shift due to the change on the refraction index due to the convection of the plasma density profile while the interaction and scattered light propagates thought it.

**B- The Effect of Density, Velocity and Temperature Gradients**

The laser intensity threshold for SBS is lowest in a homogeneous plasma (see equations. II.1.33-34 and II.1.70-71 and figs. II.2-3). As SBS is a convective instability, gradients in the electron density, temperature and expansion velocity can raise the threshold intensity significantly.\textsuperscript{25} The region where the scattered ion-acoustic and electromagnetic waves satisfy the Manley-Rowe relations and interact resonantly becomes smaller, and the propagation of wave energy out of these limited regions introduces an effective dissipation that must be overcome.

The growth rate of the SBS instability in an inhomogeneous drifting plasma can be determined using a simple heuristic approach similar to that developed in Krueer.\textsuperscript{26} A detailed and complete demonstration of these results can be found in ref. 26 (Liu, Rosenbluth and White) or ref. 2 (Liu and Kaw). From here onwards this model will be described as the convective gain or LRW (Liu-Rosenbluth-White) model for SBS in an inhomogeneous drifting plasma.

As all wavenumbers are now dependent on the position in the plasma, the wavenumber mismatch is defined as

$$\mathbf{K}(x) = \mathbf{k}_0(x) - \mathbf{k}_s(x) - \mathbf{k}_{\text{ia}}(x). \quad (\text{II.2.5})$$
At some point \( x_{\text{res}} \) inside the subcritical portion of the plasma, the wavenumber mismatch \( \mathbf{K}(x_{\text{res}}) \) is zero, so the three waves are resonantly coupled. Away from this point the mismatch grows and the resonant coupling is spoiled when a significant phase shift occurs. All plasma parameter gradients are taken along the propagation vector of the incident EM wave. The reference frame is oriented with the \( x \)-axis parallel to the propagation vector of the incident EM wave; \( \mathbf{e}_x = \mathbf{e}_k = k_0/k_0 \). The length \( L_{\text{res}} \) of the region of resonance can be estimated by using the condition on the phase shift\(^2\) \( \int_{x_{\text{res}} - L_{\text{res}}/2}^{x_{\text{res}} + L_{\text{res}}/2} \mathbf{K}(x) \cdot dx = \frac{1}{2} \). \( \text{(II.2.6)} \)

The value \( 1/2 \) for the phase shift in eqn. II.2.6 is chosen because it gives a simple result for the expression of the length of the resonance region. Taylor expanding the wavenumber mismatch around the resonance point \( x_{\text{res}} \), taken along the direction of propagation of the incident light \( \mathbf{e}_k \) gives \( \int_{x_{\text{res}} - L_{\text{res}}/2}^{x_{\text{res}} + L_{\text{res}}/2} \left( \mathbf{K}(x_{\text{res}}) + \left[ (x - x_{\text{res}}) \cdot \nabla \right] \mathbf{K}(x) \bigg|_{x = x_{\text{res}}} \right) \cdot dx = \frac{1}{2} \mathbf{e}_k \cdot \nabla \mathbf{K}(x) \bigg|_{x = x_{\text{res}}} L_{\text{res}}^2 = \frac{1}{2} . \) \( \text{(II.2.7)} \)

The gradient of the wavenumber mismatch is taken along the direction \( \mathbf{e}_k \) \( \mathbf{e}_k \cdot \nabla \mathbf{K}(x) = \mathbf{K}'(x) = \frac{\partial}{\partial x} \mathbf{K}(x) . \) \( \text{(II.2.8)} \)

Thus the length of the resonance region is thus \( L_{\text{res}} = \frac{1}{\sqrt{\mathbf{K}'(x_{\text{res}})}} . \) \( \text{(II.2.9)} \)

The effective damping rate \( \nu_{\text{ia}} \) of the ion-acoustic wave is determined by the time \( \tau_{\text{ia}} \) it takes for the wave to propagate— in the direction of increasing density gradient— out of the resonance region. The effective damping rate \( \nu_{\text{ia}} \) (see fig. II.5) can be estimated as \( \nu_{\text{ia}} = 1/\tau_{\text{ia}} = \frac{v_{\text{glA}}}{L_{\text{res}}} = \frac{v_{\text{glA}}}{L_{\text{res}}} \sqrt{\mathbf{K}'(x_{\text{res}})} . \) \( \text{(II.2.10)} \)

where \( v_{\text{glA}} = c_s \) is the group velocity of the ion acoustic wave. Similarly, the effective damping rate of the scattered EM wave \( \nu_s \)— propagating in the direction of decreasing density gradient (see fig. II.5)—is
\[ v_x = \sqrt{\frac{\tau_s}{\tau_A}} = \frac{v_{gF-M}}{L_{res}} = v_{gF-M} \sqrt{K'(x_{res})}, \] (II.2.11)

where \( v_{gF-M} = c \sqrt{1 - \frac{n_e}{n_0} \frac{\omega_s}{\omega_0}} = c \sqrt{1 - \frac{n_e}{n_0} \frac{\omega_s}{\omega_0}} \) (in SBS \( \omega_s \approx \omega_0 \), so the critical densities are approximately the same for both incident and scattered frequencies) is the group velocity for the scattered EM wave.

Inserting these effective damping rates into equation II.1.43, the Rosenbluth criterion for the \( \exp(2\pi) \) amplification of SBS in an inhomogeneous drifting plasma is obtained:

![Region of resonance for SBS in an inhomogeneous drifting plasma.](image)

Fig. II.5- Region of resonance for SBS in an inhomogeneous drifting plasma.

The region of resonance for SBS in an inhomogeneous drifting plasma is where the phase mismatch \( K = k_0 - k_s - k_{ia} = 0 \). Elsewhere SBS does not appear. Both the scattered EM (propagating up the density gradient) and the ion-acoustic (propagating down the density gradient) waves lose energy through the boundaries of this region of resonance at a rate inversely proportional to the time the waves take to traverse it.

\[ G_{SBS} = \frac{\gamma_{res}^2}{K' v_{gF-M} v_{gIA}} > 1. \] (II.2.12)

The \( \exp(2\pi) \) criterion determines the intensity threshold for instability in a inhomogeneous, drifting plasma. \( G_{SBS} \) is called the SBS amplification factor or convective gain factor. The total gain is given by

\[ G_{tot} = e^{G_{SBS}}. \] (II.2.13)
\( \gamma_{\text{res}} \) is the value of the local SBS instability growth rate at the point of resonance, \( x_{\text{res}} \), considering a homogeneous plasma. Using the isothermal model (eqn. II.1.33-35), \( \gamma_{\text{res}} \) is

\[
\gamma_{\text{res}} = \text{Im}[\omega_{\text{ia}}] = \frac{\sqrt{3}}{6} \left( \frac{k_0^2 v_{\text{os}}^2 \cos^2 \alpha}{2 \omega_0} \sum_{i=1}^{N} \omega_{\text{pi}}^2 \right)^{1/3}.
\] (II.2.14a)

For the non-local SE model (eqn. II.1.70-72), the local SBS instability growth rate at resonance, \( \gamma_{\text{res}} \), is

\[
\gamma_{\text{res}}^2 = \frac{k_{\text{ia}} v_{\text{os}}^2 \omega_{\text{pi}}^2}{8 \omega_0 |c_s - u_{\text{exp}}|} \left[ 1 + \frac{u_{\text{th},c}^2}{v_{\text{os}}^2} \right] \left( \frac{k_{\text{th}}}{k_0^2} \right)^{1/2} \cos \alpha.
\] (II.2.14b)

For stimulated amplification of the Brillouin instability, the growth rate \( \gamma_{\text{res}} \) must also satisfy the dissipative condition given by the inequality II.1.58.

To calculate the gradient of the wavenumber mismatch the gradients of the three interacting waves must be obtained

\[
K'(x) = \mathbf{\hat{e}}_k \cdot \nabla K(x) = k_0' - k_s' - k_{\text{ia}}'.
\] (II.2.15)

The ion-acoustic dispersion relationship for a drifting plasma is

\[
\omega_{\text{ia}} = k_{\text{ia}} c_s - k_{\text{ia}} \cdot \mathbf{u}.
\] (II.2.16)

The frequency of the ion-acoustic wave is constant in space. Using the angles defined in fig. II.1 to find the value of the wavenumber components along \( \mathbf{\hat{e}}_k \), and the definition II.2.8, the spatial derivative of II.2.16 becomes

\[
\omega_{\text{ia}} = 0 = k_{\text{ia}} \sin(\beta/2)c_s' - k_{\text{ia}} u' + k_{\text{ia}} c_s \sin(\beta/2) - k_{\text{ia}} u =
\]

\[
= 2 k_0 \sin^2(\beta/2)c_s' - 2 k_0 \sin(\beta/2)u' + k_{\text{ia}} c_s \sin(\beta/2) - k_{\text{ia}} u.
\] (II.2.17)

The scalelength of a quantity \( Q \) is defined as

\[
L_Q = \left| \frac{Q}{Q'} \right|.
\] (II.2.18)

Solving eqn. II.2.17 for \( k_{\text{ia}}' \) and applying the definition of scalelength II.2.18 to the gradients of the ion-acoustic speed and expansion speed, the gradient of the ion-acoustic wavenumber is
\[ k_{\alpha} = 2k_0 \frac{L_v^{-1} \sin(\beta/2) u / c_s - L_c^{-1} \sin^2(\beta/2)}{u / c_s - \sin(\beta/2)}. \] (II.2.19)

The scalelength \( L_c \) associated with the ion-acoustic speed \( c_s \), and it can be further expanded in a more extended form containing terms in electron and ion temperature, and density scalelengths.

For the incident and scattered EM wavenumbers we proceed in a similar fashion to that of the acoustic wave. Starting from the EM wave dispersion relationship

\[ \omega_0^2 = k_0^2 c^2 + \omega_{pc}^2, \] (II.2.20)

and taking the gradient gives

\[ 0 = 2k_0 k_0' c^2 + \frac{4\pi n_e e^2}{m_e} \frac{n_e'}{n_e} = 2k_0 k_0' c^2 + \frac{c^2}{1 - n_e / n_{cr}} \frac{n_e}{n_{cr}} L_n^{-1}. \] (II.2.21)

or

\[ k_0' = k_0 \frac{n_e / n_{cr}}{2(1 - n_e / n_{cr})} L_n^{-1}. \] (II.2.22)

Applying the same procedure for the scattered EM wave, a similar result follows. As

\[ \omega_s^2 = k_s^2 c^2 + \omega_{pc}^2, \] (II.2.23)

the gradient is given by

\[ 0 = 2k_s k_s' c^2 + \frac{4\pi n_e e^2}{m_e} \frac{n_e'}{n_e} = -2k_0 \cos(\beta/2) k_s' c^2 + \frac{c^2}{1 - n_e / n_{cr}} \frac{n_e}{n_{cr}} L_n^{-1}. \] (II.2.24)

The wavenumber gradients for the incident and the scattered EM waves are the same except for a geometrical factor related to the EM wave scattering angle \( \beta \) (see fig. II.1)

\[ k_s' = -k_0 \frac{n_e / n_{cr}}{2(1 - n_e / n_{cr})} \frac{L_n^{-1}}{\cos(\beta)}. \] (II.2.25)

Adding all of the wavenumber gradients, the wavenumber mismatch gradient becomes

\[ K'(x) = -2k_0 \left[ \frac{L_n^{-1} n_e / n_{cr}}{4(1 - n_e / n_{cr})} \sin^{-1}(\beta) + \frac{L_v^{-1} \sin(\beta/2) u / c_s - L_c^{-1} \sin^2(\beta/2)}{u / c_s - \sin(\beta/2)} \right]. \] (II.2.26)
Replacing eqn. II.2.25 in the threshold condition II.2.12, the equation of the intensity threshold \( I_{\text{SBS}}^\text{conv} \) for SBS convective gain in a drifting inhomogeneous, non-isothermal plasma is

\[
I_{\text{SBS}}^\text{conv} > n_\sigma m_e \frac{c_s^2}{\omega_p^2} k_0 \left[ L_n^{-1} \frac{n_e}{n_\sigma} \left[ 1 - \cos^{-1}(\beta) \right] + 4 \left( 1 - \frac{n_e}{n_\sigma} \right) \frac{L_v^{-1} \sin(\beta/2)}{c_s} \frac{u}{c_s - \sin(\beta/2)} \right]
\]

or in practical units

\[
I_{\text{SBS}}^\text{conv} > 2.5 \cdot 10^{15} \frac{Z_i T_e eV}{A_i Z_i \lambda_{\text{L,\mu m}}} \times \left[ L_n^{-1} \frac{n_e}{n_\sigma} \left[ 1 - \cos^{-1}(\beta) \right] + 4 \left( 1 - \frac{n_e}{n_\sigma} \right) \frac{L_v^{-1} \sin(\beta/2)}{c_s} \frac{u}{c_s - \sin(\beta/2)} \right]
\]

\[\text{(II.2.27)}\]

**Fig. II.6-** Convective SBS threshold in an inhomogeneous, fully ionized Silicon plasma.

The final expression II.2.27-28 for the convective SBS threshold involves a very complex relationship among the density, velocity and ion-acoustic speed scalelengths, the angle of backscattering, and the local density and expansion speed of the plasma. Usually, small scalelength gradients in electron density and velocity hinder the growth of the SBS instability, with the (ordinarily small) ion-sound speed gradient favoring the appearance of
SBS. Figure II.6 shows a plot of the convective SBS thresholds for a fully ionized Silicon plasma as a function of the density scalelength, for the cases of backscattering (β= 180°), and sidescattering (β= 110°), at two different electron temperatures.

II.3 - Finite-Length Pulse Effects

The previous theory has been developed under the assumption that the duration of laser illumination of the plasma is infinite and the response is in steady state. In practice most SBS experiments are done with finite duration laser pulses, of temporal width τ_L (usually the Full-Width, Half-Maximum [FWHM] of a Gaussian shaped pulse). In this experiment, ultrashort laser pulses of FWHM ≈ 1.6 ps are used. This introduces some additional constraints on the SBS intensity thresholds of the backscattered light.

A- SBS Saturation time

For laser light of semi-infinite duration, with a given intensity I_L (above the SBS convective threshold, eqn. II.2.27-28), incident on a finite size plasma of scalelength L_res, it has been shown in the previous section that the Brillouin instability growth is limited to a total gain

\[ G_{tot} = e^{G_{SBS}(I_L, L_{res})}, \]  

(II.2.13)

because of the dissipation through the convection of electromagnetic and ion-acoustic energy out of the region of stimulated resonance. When the total gain G_{tot} has been reached, it is said that the response of SBS to the semiinfinite duration incident laser illumination is in steady state, or convectively saturated. Increasing the laser intensity I_L exponentially increases the magnitude of total gain G_{tot}, as the convective gain G_{tot} increases linearly with the incident light intensity.

However, there is a limit to the magnitude of the total gain G_{tot} that can be reached by increasing laser intensities. At laser intensities I_L much higher than the convective threshold II.2.27, the effects of non-linear wave-particle dissipative processes, as ion-trapping and wave-breaking,28,29,30 start to be of importance. They are the main limitations to further growth of the SBS instability. When these non-linear mechanisms are strong enough to stop the growth of the total gain G_{tot} with the increase of laser intensity I_L, it is said that the absolute saturation of SBS has been reached.
Fig. II.7- SBS response to a short duration laser pulse incident on a finite homogeneous plasma.

A finite region of a plasma—in which the Manley-Rowe conditions $l_1$ for SBS resonance are satisfied—is illuminated by an incident laser pulse (iEM) a time too short to reach the saturation of the SBS response of the resulting ion-acoustic (IAW) and scattered light (sEM) amplitudes. The sequence of events illustrated in this schematic is described in the main text.

For laser light of finite duration $\tau_L$, and with a given intensity $I_L$ (above the SBS threshold), incident on a finite size plasma of scalelength $L_{\text{res}}$, the convective saturation of SBS may not necessarily be reached. If the pulse duration is too short, the Brillouin instability may not grow to the total gain level $G_{\text{tot}}$, as the SBS resonant region of the
plasma is not illuminated for a long enough time to feedback the amplification of the ion-acoustic wave.

Figure II.7 shows a schematic sequence of the events in which the incidence of a short laser pulse does not saturate the SBS response of the finite plasma. The following paragraph numbers explain the events in the correspondingly numbered parts of fig. II.7.

- 1- The incident short laser pulse (iEM) enters the finite region of SBS resonance and it interacts with a small amplitude ion-acoustic fluctuation (IAW) previously present inside this region of the plasma.

- 2- Resulting from the Bragg-like reflection of the incident laser pulse (iEM), a scattered electromagnetic wave (sEM) appears (see section I.1). The ponderomotive force associated to the combination of the incident and scattered EM waves amplifies the original ion-acoustic fluctuation (IAW). A positive feedback loop is established between increased EM scattering and ion-acoustic wave amplitude. Both the scattered EM wave (sEM) and ion-acoustic wave (IAW) grow exponentially while propagating in the plasma.

- 3- The convection of the scattered EM wave (sEM) and ion-acoustic wave (IAW) dissipates the energy of both waves out of the region of resonance. The end of the transit of the incident laser pulse (iEM) out of the region of resonance stops the growth of SBS before it reaches saturation.

- 4- The faster scattered EM wave (sEM) leaves the region of resonance before that the ion-acoustic wave (IAW). The amplitudes of the scattered EM wave (sEM) and ion-acoustic wave (IAW) remains unchanged, but are lower than if they have reached saturation.

If the laser pulse lasts long enough to saturate the convective SBS response, the remainder of time in which the laser pulse is present does not further increase the plasma reflectivity and scattered light intensity. So a laser pulse of finite duration, but longer in width than the time needed for the plasma to reach response saturation, will behave in the same way than a semiinfinite laser pulse which starts at a definite time but never switches off. Figure II.8 shows a schematic sequence of the events leading to the convective saturation of the SBS response by an incident long laser pulse. The following paragraph numbers explain the events in the correspondingly numbered parts of fig. II.8.
Fig. II.8 - SBS response to a long duration laser pulse incident on a finite homogeneous plasma.

A finite region of a plasma—in which the Manley-Rowe conditions for SBS resonance are satisfied—is illuminated by an incident laser pulse (iEM) a time long enough to reach the saturation of the SBS response of the resulting ion-acoustic (IAW) and scattered light (SEM) amplitudes. The sequence of events illustrated in this schematic is described in the main text.

- 1- The incident long laser pulse (iEM) enters the finite region of SBS resonance and it interacts with a small amplitude ion-acoustic fluctuation (IAW) present inside this region of the plasma.
• 2- Resulting from the Bragg-like reflection of the incident laser pulse (iEM), a scattered electromagnetic wave (sEM) appears (see section I.1). The ponderomotive force associated to the combination of the incident and scattered EM waves amplifies the original ion-acoustic fluctuation (IAW). A positive feedback loop is established between increased EM scattering and ion-acoustic wave amplitude.

• 3- Both the scattered EM wave (sEM) and ion-acoustic wave (IAW) grow exponentially while propagating in the plasma. The convection of the scattered EM wave (sEM) and ion-acoustic wave (IAW) out of the region of resonance is compensated by the infusion of additional energy by the incident laser pulse (iEM).

• 4- The convection of the scattered EM wave (sEM) and ion-acoustic wave (IAW) dissipates the energy of both waves outside the region of resonance at the same rate that it is replenished by the incident laser pulse (iEM). The amplitudes of the scattered EM wave (sEM) and ion-acoustic wave (IAW) reach their maxima. The convective saturation of the SBS response is reached. The amplitudes of both waves remain at the saturation level until the incident laser pulse (iEM) has propagated entirely out of the region of SBS resonance.

Following the treatment of Giacone, McKinstrie, and Betti \(^{31}\) for the SBS response of a stationary, homogeneous, spatially finite plasma of length \(L_{\text{res}}\) to a laser pulse of semi-infinite duration, a rough estimation of the minimum laser pulse temporal width needed for SBS to convectively saturate. The spatiotemporal evolution of the amplitude of the scattered electromagnetic wave is

\[
E_s(x,t) = v_{\text{gEM}} \int_0^t G_{ss}(x,t') \, dt',
\]  

(II.3.1)

where the Green’s function

\[
G_{ss}(x,t) = \frac{\gamma}{v_{\text{gEM}} + v_{\text{glA}}} \frac{x + v_{\text{glA}} t}{v_{\text{gEM}} t - x} I_1 \left[ \frac{\gamma}{v_{\text{gEM}} + v_{\text{glA}}} \sqrt{(v_{\text{gEM}} t - x)(x + v_{\text{glA}} t)} \right] \times 
\exp \left[ - \frac{\nu_{\text{col}} (v_{\text{gEM}} t - x) + \nu_{\text{Lan}} (x + v_{\text{glA}} t)}{v_{\text{gEM}} + v_{\text{glA}}} \right] H(v_{\text{gEM}} t - x) H(x + v_{\text{glA}} t) + 
H(x + v_{\text{glA}} t) \delta(v_{\text{gEM}} t - x) \exp(-\nu_{\text{col}} t). 
\]  

(II.3.2)

The Green’s function describes the effect on the scattered electromagnetic wave response at the coordinates \((x,t)\) of the impulse provided by the incident electromagnetic
wave at the coordinates (0,0). Figure II.9 shows plot for the SBS impulse response (eqn II.3.2) and its integral (eqn. II.3.1) for a Silicon plasma with parameters are $T_e = 1$ keV, $n_e/n_c = 0.1$ and $L_{res} = 100$ μm, illuminated by a semi-infinite laser pulse of intensity $I_c = 10^{15}$ W/cm². Under the low laser intensity regime condition

$$\frac{2\gamma^2 L_{res}}{\left(\frac{\nu_{col}}{v_{gEM}} + \frac{\nu_{Lan}}{v_{gIA}}\right)v_{gEM} v_{gIA}} \ll 1,$$

an asymptotic approximation of the Green's function can be written as

$$G_{ss}(x,t) = \gamma^2 \frac{x + v_{gIA} t}{(v_{gEM} + v_{gIA})^2} \exp\left[-\frac{\nu_{col}(v_{gEM} t - x) + \nu_{Lan}(x + v_{gIA} t)}{v_{gEM} + v_{gIA}}\right].$$

$G_{ss}$ decays monotonically in time, independently of the interaction laser intensity. For a plasma of finite length $L_{res}$, the value of the integral II.3.1, taken between times 0 and $t$, gets to 95% of its maximum at $t = \tau_{low}$, is reached, where

$$\tau_{low} = \frac{L_{res}}{v_{gEM}} \left[1 + \frac{3}{v_{gIA} L_{res}} \frac{\nu_{col}}{v_{gEM}} + \frac{\nu_{Lan}}{v_{gIA}}\right].$$

The integral II.3.1 remains nearly constant when integrated for upper time limits with $t > \tau_{low}$. The saturation of the SBS convective response in this regime is then reached.

The theory of Giacone et al. applies to a stationary, homogeneous finite plasma. As the phase speed $v_{gEM}$ (c) of the incident laser pulse is much larger than the fastest portion of an expanding plasma ($v_{exp} \approx 10^{-4} c$), the stationary condition can be dropped easily. The inhomogeneity in density can be ignored by requiring that the three-wave phase mismatch $K' (x_{res})$ within the finite length SBS resonance region to be small (see eqn. II.2.7).

For an expanding inhomogeneous plasma, the finite length of the region of resonance is $L_{res} \approx \frac{1}{\sqrt{K'(x_{res})}}$ (eqn. II.2.9). Near the SBS convective gain threshold (eqn. II.2.12), the convective saturation condition in the low-intensity laser regime (II.3.3) becomes
Fig. II.9- SBS impulse response of a finite scalelength, homogeneous Silicon plasma to a semiinfinite laser pulse. The incident laser light has intensity $I_L = 10^{15}$ W/cm$^2$. The plasma parameters are $T_e = 1$ keV, $n_0/n_{cr} = 0.1$ and $L_{res} = 100$ µm. Convective saturation response is reached when $t = \tau_{low}$.

$$\frac{\gamma}{\sqrt{\frac{g_{EM}^2 g_{IA}^2}{4^4 4^4 4^4 4^3}}} \ll 1.$$  \hspace{1cm} (II.3.6)

For usual experimental conditions, the left hand side of II.3.6 is not greater than $10^{-4}$, so the low laser intensity regime condition is met. The low-intensity regime condition to achieve convective saturation of the SBS instability within the time period while the laser interaction pulse is illuminating the region of resonance (eqn. II.3.5) is
\[
\tau_{\text{low}} = \frac{L_{\text{res}}}{v_{\text{gIA}}} \left[ \frac{1}{2} \left( \frac{v_{\text{col}}}{v_{\text{gEM}}} + \frac{v_{\text{Ian}}}{v_{\text{gIA}}} \right) \sqrt{v_{\text{gEM}} v_{\text{gIA}}} \right] < \tau_L.
\]

For the actual experimental conditions, the term in the square brackets is very small (between $10^{-3}$ to $10^{-5}$), and gets smaller with longer scalelength plasmas. The condition (II.3.7) therefore requires that the interaction laser pulse must illuminate the region of resonance $L_{\text{res}}$ during a very small fraction of the time the ion-acoustic wave takes to traverse this region, $L_{\text{res}}/v_{\text{gIA}}$. The laser pulse width $\tau_L$ must be longer than the illumination time $\tau_{\text{low}}$ to reach convective saturation of the SBS instability in the low intensity regime. Rearranging factors in eqn. II.3.7, and requiring

\[
\frac{3}{4} \frac{v_{\text{gIA}}}{v_{\text{gEM}}} \frac{K'}{v_{\text{gEM}} \left( \frac{v_{\text{col}}}{v_{\text{gEM}}} + \frac{v_{\text{Ian}}}{v_{\text{gIA}}} \right)} \ll 1,
\]

the condition for the time to reach response saturation reduces to

\[
\tau_{\text{low}} = \frac{L_{\text{res}}}{v_{\text{gEM}}} < \tau_L.
\]

Therefore, SBS response saturation will be reached if the interaction laser pulse width $\tau_L$ is longer than the time the incident electromagnetic wave takes to traverse this region, $L_{\text{res}}/v_{\text{gEM}}$. For ultrashort pulses with $\tau_L$ in the picosecond range, it means that the length effective region of resonance $L_{\text{res}}$ for SBS must not be greater than 200 to 400 $\mu$m.

Batson, Baldis et al.\textsuperscript{32,33,34,35} have also analyzed the problem of the interaction of an ultrashort laser pulse with a finite homogeneous plasma in detail. For the experimental conditions of the laser pulse and plasma conditions described elsewhere in this thesis, their analysis reduces to similar results to those derived from the Giacone, McKinstrie and Betti treatment. For the SBS experiment described in this thesis, the conditions II.3.3 and II.3.7 are always satisfied.
B- The Temporal threshold

The growth period of the SBS excited ion-acoustic wave can be restricted by the short laser pulse duration, even for the most favorable plasma conditions. The laser pulse must be intense enough to allow the exponential increase of the SBS instability during the duration of its illumination of the plasma. This imposes a second intensity threshold condition to be overcome

$$\gamma \tau_L \geq 1$$  \hspace{1cm} (II.3.10)

which, applying eqn. II.1.34 and using practical units, translates into the inequality

$$I_{\text{SBS}\text[\text{temp}] \text[V/cm}^2] \geq \frac{2.8 \cdot 10^{14}}{Z_i \tau_{\text{Lps}} n_e/n_{\text{cr}} \cos(\beta/2) \sqrt{1 - n_e/n_{\text{cr}}}} \frac{1}{A_i}$$  \hspace{1cm} (II.3.11)

This temporal laser intensity threshold—with values between $10^{13}$ to $5 \cdot 10^{14}$ W/cm$^2$—is very easily attained within the usual experimental parameters, even for ultrashort laser pulses. Figure II.10 shows the temporal SBS threshold intensity $I_{\text{SBS}\text[\text{temp}]}$ for a homogeneous stationary fully-ionized Silicon plasma for normally incident 1.6 ps FWHM laser pulse for two different plasma densities.

![Graph showing SBS threshold intensity vs. electron temperature](image)

**Fig. II.10 - Temporal SBS threshold intensity for a homogeneous stationary fully-ionized Si plasma for normally incident 1.6 ps FWHM laser pulse.**
II.4 - Seeding the Instability

When the intensity of the incident laser light exceeds both the dissipative (collisional plus Landau damping) and the convective Rosenbluth thresholds of the SBS instability, a feedback mechanism linking the increase of the scattered light intensity with the amplitude of the ion density fluctuation is established. As a result, the intensity of the Brillouin scattered light $I_s$ grows exponentially from an initial electromagnetic seeding wave, of intensity $I_{s0}$. Therefore,

$$I_s = I_{s0} \left( n_e, T_e, \{ n_i, T_i \}_{i \in \{ N \}} \right) e^{G_{SBS}}.$$  \hspace{1cm} (II.4.1)

Several different processes have been considered as the initial seeding source of the Brillouin backscattered electromagnetic wave. A short summary of some of them follows.

A. Seeding by the Thermal EM Noise Background

The most common source for the seeding of SBS is the background of thermally originated electromagnetic plus electrostatic (plasma wave) random fluctuations present in any plasma.\textsuperscript{36} Even in a stable, equilibrium plasma, there are a finite level of electromagnetic and plasma waves. These are the ‘normal modes’ of the system, and represent degrees of freedom which are excited in thermal equilibrium. Plasma and EM waves are emitted by the particles moving in the plasma and absorbed continuously again (via collisional and Landau damping). The balance between emission and absorption leads to thermal noise levels of electromagnetic field and plasma density fluctuations.

The calculation\textsuperscript{37,38} of these thermal noise levels—using test charge methods in the Vlasov theory to include particle discreteness effects—is too lengthy and intricate to be described here, so only the end results will be shown. In the Brillouin case, the backscattered light flux resulting from the thermal electrostatic (ion-acoustic waves) fluctuations takes the form

$$F_s = \frac{d^2 I_s}{d\Omega d\lambda_s} = \frac{CT_L}{\lambda^4} \frac{\omega_s}{2 \left( \varphi \omega_a - \omega_s \right)} e^{G_{SBS} \left( \omega_s \right)} - 1 \right],$$  \hspace{1cm} (II.4.2)

where the units for the flux $F_s$ are intensity of radiation per unit solid angle $\Omega$ per unit vacuum wavelength. The effective temperature $T_L$ is given by
\[ T_L = \frac{T_e}{1 + Z \sqrt{\frac{T_e/m_e}{T_i/m_i}} e^{-\frac{1}{2} \left( \frac{T_e}{T_i} - 1 \right)}} \]

Similarly, the expression for the Brillouin backscattered flux seeded by the thermal electromagnetic noise is roughly

\[ F_i = \frac{d^2 F}{d\Omega d\lambda_s} = \frac{c T_e}{\lambda_s^3} e^{G_{SBS}(\omega_s)}. \]

The contribution of the electrostatic density fluctuations far exceeds that of the electromagnetic thermal noise, because the ratio \( \omega_s/2(\omega_0 - \omega_s) \approx c/4c_s \) is very large for SBS (5 \cdot 10^2 to 10^4 for usual experimental conditions). As ion-acoustic fluctuations (phonons) carry less energy than electromagnetic fluctuations (photons), they are easier to excite, so more of them appear in the plasma and the exponential buildup of the scattered light will proceed from an enhanced source.

**B- Seeding by Accelerated Ion Jets**

Another mechanism for seeding the Brillouin scattering of the incident laser light is the presence of highly energetic jets of accelerated ions,39,40,41 moving down the density gradient at supersonic speeds, and exciting an enhanced background of ion-acoustic waves through decay by streaming-like instabilities. Ion jets may be generated by resonant absorption or Langmuir wave collapse near the critical surface of the expanding plasma.

When the ion jet speed satisfies with the relationship:42

\[ |k_{ia} (u_{i,jet} - u_{exp})| < \omega p e, \]

a streaming instability is triggered, and the ion jet decays into ion-acoustic waves of wavenumber \( k_{ia} \) (see fig. II.11). The condition II.4.5 can be rewritten as

\[ \frac{u_{i,jet} - u_{exp}}{c} < \frac{1}{2} \sqrt{\frac{n_e}{n_{cr}}} \left( 1 - \frac{n_e}{n_{cr}} \right). \]
Fig. II.11- Brillouin scattering from an ion beam decaying into ion-acoustic waves through an streaming instability.

The incident laser light now interacts with the ion-acoustic waves $k_{ia}$ resulting from the decay of ion jets. When the ion jet speed is such that $|k_{ia}(u_{i,jet} - u_{exp})| < \omega_{pe}$, it triggers a streaming plasma instability.

The resulting enhanced background of ion-acoustic waves is illuminated with laser light, and two cases are possible:

- **Simple Brillouin scattering**, $k_{ia} \neq 2k_0$—The laser light Bragg-like scatters from the enhanced background of ion-acoustic fluctuations. The Manley-Rowe conditions I.1 are satisfied, but no positive feedback loop is established between the growth of the scattered light amplitude and that of the ion-acoustic fluctuations. Two observable consequences of this kind of interactions are:
  
  i- The observed Brillouin scattered spectrum appears extremely blueshifted, because most of the ion-acoustic fluctuations spread far out into the faster, less dense regions of the expanding plasma, where the condition II.4.6 is satisfied. Also, the conservation of the momentum carried by the original highly energetic ion jet implies that most of the ion-acoustic waves should counter-propagate with respect to the direction of the incident laser light, limiting even more the region in which Brillouin scattering is feasible.

  ii- The observed Brillouin scattered spectrum should show features covering a large spectral bandwidth. As the incident laser light non-resonantly Brillouin scatters
from the broadband ion-acoustic frequencies and wavenumbers present in the interaction region, the Manley-Rowe relations imply that the resulting Brillouin scattered light must correspondingly be also highly spectrally spread. Unfortunately, research on this area is still in its infancy and an expression for the spectral distribution of the ion-acoustic wave spectrum has yet to be found to make more quantitative predictions feasible.

- **Stimulated Brillouin scattering, \( k_{ia} = 2 k_\eta \)**—The laser light scatters from the small part of the fast ion-jet produced background of ion-acoustic fluctuations that not only satisfies the Manley-Rowe conditions I.1, but also allows that a resonant feedback loop to be established between the growth of the scattered light amplitude and that of the ion-acoustic fluctuations. The observable SBS spectral features have now a narrow bandwidth, but as most of the ion-acoustic fluctuations are still localized in the faster, thinner regions of the expanding plasma. They are extremely blueshifted.

Again, the lack of an equation for the spectral distribution of the ion-acoustic wave spectrum impedes to find expressions similar to II.4.2-3 for the initial background flux from which the stimulated Brillouin scattered light amplifies.

**C- The Coupling of SBS to SRS**

Another process enhancing the level of electromagnetic field and density fluctuations is the coupling between SBS and Stimulated Raman Scattering (SRS).\(^{43,44,45,46}\) Although SRS develops from a noise level lower than that of SBS by a factor \( \omega / \omega_{pe} \), it has a larger growth rate. The SRS driven electron plasma waves can originate parametric decay or modulational instabilities\(^{47}\) which lead to their cascading and/or collapse. These nonlinear processes may generate ion-acoustic wave turbulence, therefore enhancing any noise background of the density fluctuations in the plasma from which SBS could, in turn, develop. This enhancement of the ion-acoustic background results in a higher scattered light intensity, but the expression for the SBS convective threshold (eqn. II.2.27-28) remains the same.

An observable consequence of the presence of a SRS enhanced ion-acoustic background is that the Brillouin scattering must remain in the region where \( n_e < n_{cr}/4 \). Therefore, the backscattered Brillouin spectrum must appear strongly Doppler-blueshifted, because this very underdense region is typically undergoing rapid expansion (\( v_{exp} > 5 c_s \)).
However the comparison of numerical calculations to a previous experiment\textsuperscript{35} has indicated that it takes too long for a SRS enhanced background to play a significant role in the seeding of SBS when an ultrashort laser pulse is used for interaction.

\section*{II.5 - Translating SBS Theory into an Experiment}

The theoretical results derived in previous sections of this chapter must be integrated to determine the expected spectral and energy features of the backscattered light coming from the interaction of an ultrashort pulse laser with a solid planar target.

The plasma formation, expansion and hydrodynamics are described extensively in section III.1. For the current discussion, it is sufficient to know that the plasma has a density gradient in the direction of the laser propagation and is expanding supersonically with increasing velocity. Electron and ion temperatures also have small gradients, but it is a good approximation to assume the plasma is isothermal for all species.

To avoid damage to the laser system with specularly reflected light, plus to maximize the contrast of the backscattered light signal to the background optical noise, the normal of the planar solid target is at an angle with respect to the direction of propagation of the incoming light. The geometry is shown in fig. II.12.

The inclination of the solid target makes the plasma critical surface specularly reflect the incident laser beam. The specularly reflected light intensity is much lower than the incident light intensity due to collisional absorption in the bulk of the plasma and to resonant absorption close to the critical surface.

\subsection*{A- Ideal and observable spectra}

Below the SBS threshold, the backscattered spectrum only shows a single peak resulting from the Lambertian reflection of the incident interaction pulse from the critical surface, plus specular reflection from the optics of the laser system and diagnostics optics. This peak shows a small blueshift because the critical surface is usually moving towards the observer.

Ideally, when the incident laser intensity is increased above the SBS convective threshold value and neglecting the Doppler shift resulting from the plasma expansion a second redshifted feature appears (see the upper left plot in fig. II.13). This feature is the
spectral signature of Brillouin scattering, and it increases its amplitude exponentially with increasing incident intensity. But as the plasma is expanding supersonically outwards, the observed Brillouin peak is strongly Doppler-shifted to the blue side of the Lambert-scattered peak (see the upper right, light framed plot in fig. II.13). Measured spectra with single SBS features are shown in section IV.1B for silicon, section IV.2B for Gold, and section IV.3B for CH. Analysis of these spectra is found in section IV.4B.

Fig. II.12- Observable stimulated Brillouin light scattering by a plasma expanding from a plane solid target.

The observable SBS light comes from two different scattering events:
A- Brillouin backscattered light results from the interaction of the incident laser pulse with the expanding plasma.
B- Brillouin sidescattered light comes from the interaction of the laser pulse reflected by the critical density surface with the expanding plasma.

A further increase of the incident interaction laser intensity, results in the specularly reflected intensity from the critical surface being high enough to overcome the Brillouin sidescattering threshold for twice the angle of incidence of the light to the target normal. Now two redshifted Brillouin features (neglecting the plasma expansion Doppler effect) are expected to appear in the spectrum (see the lower left plot in fig. II.13). The higher peak results from Brillouin backscattering of the incident light, while the redder and smaller
Fig. - II.13- Appearance of the observable SBS spectrum.

**Left column:** In a stationary plasma, a Brillouin peak appears to the red side of the backreflected laser peak (wavelength - \( \lambda_L \)) when a high-intensity laser beam interacts with the plasma expanding from a solid planar target at an angle with respect to the incident beam direction (upper left plot). With increased interaction intensity, an additional redder SBS feature also appears (lower left plot). The more redshifted peak is associated with the backscattering interaction of the incident laser light. The second, usually smaller peak results from the sidescattering interaction of the critical surface specularly reflected laser light.

**Right column:** The highly supersonic expansion of the plasma shifts both peaks to the blue side of the laser peak. The upper right plot (light frame) shows the observed spectrum when the laser intensity is high enough to see only Brillouin backscattering. The lower right plot (heavy frame) shows the observed spectrum when the laser intensity is high enough to see both Brillouin backscattering and sidescattering.
comes from Brillouin sidescattering of the specularly reflected light of the critical surface as it passes out of the plasma. With increased intensity both features grow in amplitude.

However, the expansion of the plasma again strongly Doppler-shifts the observed Brillouin features to the blue side of the Lambert-reflected peak position (see the lower right, heavy framed plot in fig. II.13). The measured spectra with double SBS features are shown in section IV.1B for Silicon, section IV.2B for Gold, and section IV.3B for CH. Analysis of these spectra is found in section IV.4B.

B- Measuring the SBS threshold

The total (for all the wavelength range observed) and partial (for a particular feature wavelength) intensities of the backscattered light can be measured, integrated over the time the SBS instability is active. Plotting the strength of the measured signal versus the interaction laser intensity (for the same preformed plasma parameters) a S-shaped curve like that shown in fig. II.14 is expected to appear. Three regions are noticeable:

- **No SBS:** The interaction laser intensity is below the SBS threshold. Here the scattered signal increases linearly with the laser intensity, so the reflectivity is constant and small (\( R \sim 10^{-5} \)). The reflected signal comes mainly from Lambertian reflection on the critical surface of the expanding plasma plus spurious reflections from the surfaces of the optics of the laser system and collection optics of the diagnostics.

- **Linear SBS:** Above a certain value of the interaction laser intensity—the SBS threshold, the reflectivity of the plasma experiences a sudden jump. From this point on the reflectivity increases exponentially (up to absolute reflectivity values of \( R \sim .1 \) to \(.3 \)) until it reaches saturation. The SBS parameters are estimated from the linear theory developed on sections II.1 to II.3.

- **Saturated SBS:** Above the so-called saturation intensity, the reflectivity of the plasma stops growing and flattens out (remaining at \( R \sim .1 \) to \(.3 \)). Non-linear dissipative processes in the plasma absorb the increasing laser intensity instead of scattering it. The detected signal still grows linearly with the interaction laser intensity, but this increase comes from the same constant reflectivity sources (optical noise plus Lambertian reflection at target) that are apparent in the region below threshold.
Fig. II.14 - The ideal behavior of the total SBS signal with increasing laser interaction intensity.

Three different regions are expected to be seen. (1) No SBS: where the backscattered signal increases linearly with the interaction intensity, and results mainly from Lambertian reflection from the plasma critical surface (plus spurious reflections from the diagnostics light collection optics). (2) Linear SBS: where, after surpassing the SBS threshold, the scattered signal grows exponentially from the background level of ion-acoustic noise. (3) Saturated SBS: where the plasma cannot support additional growth of the SBS signal, so non-linear damping mechanisms overcome the SBS process.

From measurement to measurement during the experiment, the parameters of the plasma cannot be kept fixed to a certain set of density, expansion speed and temperature values, but drift around the expected ones. The measurement consists of the superposition of several S-shaped curves as that shown in fig. II.15, each associated to a particular set of initial plasma conditions.

Both the SBS intensity threshold and the saturation threshold values have some variation with the plasma formation intensity. This effect gives a level of uncertainty to the position of these thresholds and makes it difficult to estimate of the value of the growth rate, due to the resulting scatter of discrete measured points pairs (laser intensity, backscattered signal). Plots of the measured SBS backscattered reflectivities are shown in § IV.1A for Silicon, § IV.2A for Gold, and § IV.3A for CH. Analysis of these energy measurements is found in section IV.4A.
Fig. II.15 - Multiple SBS growth curves appear on the actual measurement of the backscattered energy.
Changing conditions on the plasma from measurement to measurement result in the appearance of multiple SBS growth curves, blurring the boundaries of instability threshold and saturation.
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III- EXPERIMENTAL SET-UP

IN THIS CHAPTER, the methods and equipment used to perform the experiment are presented. The formation of a plasma by ultrashort laser pulses with solid planar targets plus its subsequent isothermal expansion is analyzed theoretically, numerically and experimentally. The TableTop Terawatt (TTT) Chirped Pulsed Amplification laser system is described. The pulse splitting and stacking using a Michelson interferometer is explained. A description of the generation of high contrast laser pulses using of a saturable absorber is given. The Zeta vacuum chamber and its targets are described. The spectral and energy diagnostics for both incident and backscattered light are described, including their calibration. The measurement of the preformed plasma scalelength by the use of an integrating sphere by the laser light absorbed by the plasma is shown.

III.1 - PLASMA FORMATION BY ULTRASHORT LASER PULSES

Creating the plasma conditions required for SBS to be above the convective threshold using only a single intense ultrashort laser pulse is extremely difficult. During the interaction of an ultra-short (~1 ps), intense laser pulse (>10^{15} W/cm^2) with a solid target, the plasma is formed with a density higher than critical, and undergoes limited hydrodynamic expansion. Its density scalelength L_n is smaller than the laser wavelength \lambda_L (\lambda_L = 1054 nm).^{1}

A plasma formed by an intense ultrashort laser pulse must be allowed to expand for a time which is sufficiently long to ensure that the underdense plasma density scalelength is tens of laser wavelengths \lambda_L. The convective threshold intensity for SBS scales as 1/L_n^{2-3}.^{2-3}

Using two ultra-short laser pulses separated in time reduces the difficulties of the single pulse case, so the SBS threshold is lower. Figure III.1 shows the sequence of events conducing to plasma formation, expansion and SBS interaction through the use of two temporally separated pulses. The first ultrashort pulse primarily heats the target electrons through collisional and resonant absorption of the EM wave. Subsequently, the electron thermal energy is shared collisionally with the ions. The electron fluid acts as a heat reservoir from which the ions draw thermal energy.
$t < 0$: 
The plasma formation pulse enters the target chamber, 
$\lambda_L = 1053 \text{ nm}, \text{ FWHM} = 1.6 \text{ ps}$, 
$10^{-8}$ pedestal/peak contrast, 
$10^{14} \text{ W/cm}^2 \leq I_{\text{for}} \leq 5 \times 10^{16} \text{ W/cm}^2$.

$t = 0$: 
Plasma formation pulse hits the target and creates a very short scalelength, denser than critical plasma, $L_n = 0.2 \mu \text{m}$.

$0 < t < \Delta \tau$: 
The preformed plasma expands isothermally during a time $\Delta \tau$, 
$0 \leq \Delta \tau \leq 1500 \text{ ps, } 25 \mu \text{m} \leq L_n \leq 250 \mu \text{m}$. 
The interaction pulse enters the target chamber, $\lambda_L = 1053 \text{ nm, } \text{ FWHM} = 1.6 \text{ ps}$, 
$10^{-8}$ pedestal/peak, $0.1 \leq I_{\text{int}}/I_{\text{for}} \leq 10$, 
$10^{14} \text{ W/cm}^2 \leq I_{\text{int}} \leq 5 \times 10^{16} \text{ W/cm}^2$.

$t = \Delta \tau$: 
The interaction pulse hits the expanded subcritical plasma. If $I_{\text{int}}$ is greater than the threshold, backscattered SBS is generated. 
The interaction beam reflected from the critical surface produces sidescattered SBS if its intensity is over the threshold.

Fig. III.1 - Plasma formation and SBS interaction sequence.

The underdense part of the plasma expands isothermally, due to the high electron heat conductivity and the lack of additional heating by the laser. The second delayed ultrashort pulse interacts with an isothermal expanding plasma with an exponential density profile given by$^4$

$$n(x, \Delta \tau) = n_{cr} e^{-(x - v_{cr} \Delta \tau)/c \cdot \Delta \tau}$$

$$v(x, \Delta \tau) = x/\Delta \tau$$  \hspace{1cm} (III.1.1)

where $n_{cr}$ is the plasma critical density (see eqn. I.4), $v_{cr}$ is the speed of the plasma's critical surface towards the observer, $x$ is the position along the plasma line of expansion,
and $\Delta \tau$ is the delay between the preforming and the interaction pulses. This is described in more detail in subsection III.1B. Therefore the isothermal plasma density and velocity scalelengths are given by

$$L_n = c_s \Delta \tau$$
$$L_v = x.$$  \hspace{1cm} (III.1.2)

Numerical simulations using the 1-D Lagrangian code LILAC (see § III.1D at page 80) substantiate this simple model. The plasma expansion speed is comparable to or higher than the plasma's ion-sound speed ($>10^6$ cm/s) in its underdense regions.

LILAC simulations predict that the ponderomotive force associated to the plasma forming prepulse compresses the target surface to densities several times higher than critical in a region smaller than a laser wavelength. Therefore, SBS is extremely difficult to be excited for single ultrashort pulse interactions, with convective threshold intensities in the order of $10^{19}$ to $10^{24}$ W/cm$^2$. A delay period for expansion of the plasma is needed, followed by a second interaction pulse, if SBS is to be detected.

**A- Plasma Formation by Light Absorption**

A powerful ultrashort (in the ~0.1-1 ps regime) high-contrast laser pulse interacting with solid target produces a plasma with a density higher than critical, and with very energetic electrons.\textsuperscript{5} Most of the absorbed laser energy is transferred to the electrons, and the ions remain mostly cold, because the electron-ion the collision time is much longer than the laser pulse duration. However, this plasma does not maintain a step-like ion density profile but experiences a limited expansion of its ion component because the ambipolar electric field arising from charge separation at the target surface.\textsuperscript{6,7,8} Fast energetic electrons leaving the plasma create this field. A denser than critical plasma, with a density profile steeper than that resulting from isothermal expansion and a density scalelength smaller than the laser wavelength ($L_n < \lambda_L$) is created. Two energy absorption processes dominate the formation and heating of a plasma by ultrashort laser light, collisional absorption and resonant absorption.

**a- Collisional Absorption**

In collisional absorption, the electrons gain energy from the laser beam by thermalizing the oscillatory energy of the electromagnetic field through collisions with the colder ion background. This process is independent of the laser polarization, but dependent on the
angle of incidence $\phi$ of the beam with respect to the plasma electron density gradient and the profile of this gradient $n_e = n_e(x, L_n)$, due to the location of the turning point and the length of the laser path inside the plasma. Figure III.2 shows a schematic of the beam-plasma geometry for collisional absorption.

Along a plasma slab—with a monotonously decreasing density gradient, $\hat{e}_x \cdot \nabla n_e < 0.$ in the subcritical region, $x > 0$—, the electromagnetic wave locally transfers energy to the electrons at a collisional damping rate\(^9\)

$$u_{\text{coll}} = u_{\text{ei}} \frac{\omega_{\text{pe}}^2}{\omega_0^2} = u_{\text{ei}}(x) \frac{n_e(x)}{n_{\text{cr}}}.$$  \hfill (III.1.3)

![Diagram of laser beam absorption](image)

**Fig. III.2 - Absorption of an obliquely incident laser beam in an inhomogeneous plasma.**

The plasma is an inhomogeneous stratified slab with its physical parameters changing only in the $x$-direction. The beam plane of incidence is defined by the beam propagation vector and the vector antiparallel to the plasma density gradient. Collisional absorption is independant of the polarization of the incident beam, but resonant absorption only appears when there is a non-zero P-polarization component. P-polarization beams have their electric field vector oscillating parallel to the plane of the page. S-polarization beams have their electric field vector oscillating in and out of the page.

As the beam passes through the plasma its initial (vacuum) intensity $I_{L0}$ decreases monotonically. When leaving the plasma, the laser beam has an intensity $I_{Lf}$. The collisional fractional absorption in the plasma, for a laser beam incident at a plasma with an angle $\phi$, is defined as...
\[ f_A^c = 1 - \frac{I_{lf}(\phi, L_n, v_{ei})}{I_{l0}}. \] (III.1.4)

The value of the collisional fractional absorption is sensitive not only to the laser beam angle of incidence, but also to the shape of the electron density profile. A rather involved calculation using WKB theory\(^{10}\) allows the integration of the infinitesimal contributions to the collisional damping of the electromagnetic wave along its path inside the plasma. For a linear density profile, \(n_e = n_{cr} (1 - x/L_n)\), the collisional fractional absorption is

\[ f_A^c = 1 - \exp \left( -\frac{32}{15} \frac{v_{ei} L_n}{c} \cos^5 \phi \right). \] (III.1.5)

**Silicon, \(T_e = 1000\) eV**

![Graph showing collisional fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a 1 keV Silicon plasma.]

Fig. III.3 - Collisional fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a 1 keV Silicon plasma.
CHAPTER III - EXPERIMENTAL SET-UP

The electron-ion collisions frequency $\nu_{ei}^*$ is the same as in eqn. II.1.38, but now evaluated at the critical electron density,

$$
\nu_{ei}^* = \nu_{ei}(n_e = n_{cr}) = \begin{cases} 
\frac{1}{3(2\pi)^{3/2}} \frac{\omega_0^3}{n_{cr} u_{th,e}} \sum_{i=1}^{N} \frac{n_i(n_e=n_{cr})}{n_{cr}} Z_i^2 \ln \Lambda_{Ci} & , \text{for multiple species.} \\
\frac{1}{3(2\pi)^{3/2}} \frac{\omega_0^3}{n_{cr} u_{th,e}} Z_i \ln \Lambda_{Ci} & , \text{for a single species.}
\end{cases}
$$

(III.1.6)

The weak dependence of $\nu_{ei}^*$ of the Coulomb logarithm $\ln \Lambda_{Ci}$ (eqn. II.1.45) on the electron density has been neglected. Figure III.3 shows a plot of the collisional fractional absorption $f_A^C$ at different electron density scalelengths $L_n$ for a 1054 nm laser beam interacting with a 1 keV Silicon plasma. It can be seen that collisional absorption is important mostly in longer density scalelength plasmas ($L_n > 10 \lambda_L$) and it reaches its maximum value for normal incidence ($\phi = 0$).

\textit{b- Resonant Absorption}

When P-polarized light is obliquely incident on an inhomogeneous plasma slab, the longitudinal electric field component (along the local density gradient of the plasma) drives a charge oscillation in the plasma, which produces resonant electrostatic plasma waves in the region close to the critical surface. The beam-plasma geometry for resonant absorption is shown in fig. III.2; p-polarized light has its electric field vector coplanar to the plane of incidence of the laser beam.

Subsequently the beam reaches the turning point ($x = L_t$) near the plasma critical surface ($x = 0$), it reflects and propagates out of the plasma, but its oscillating electric field tunnels into the critical density region to excite an electron plasma wave at a frequency $\omega_{pe} = \omega_0$. The electric field vector has a component parallel to the direction of the density gradient, $\mathbf{E} \cdot \nabla n_e \neq 0$. Part of the incident wave energy is transferred to the resonant electron plasma wave, which dissipates it through collisional and Landau damping, wave-particle interactions, or by convection outside the plasma critical region. The incident beam intensity $I_{L,0}$ decreases and after leaving the near-critical region of the plasma, the laser beam has an intensity $I_{L,f}$. The fractional resonant absorption of a plasma, is defined as

$$
f_A^R = 1 - \frac{I_{L,f}(\phi, L_n)}{I_{L,0}}.
$$

(III.1.7)
depending only on the laser beam angle of incidence $\phi$ and density profile. It is independent of the material composition of the plasma. It must be remembered that resonant absorption does not appear in homogeneous plasmas or for S-polarized light.

The detailed calculation of the fractional resonant absorption is rather involved and it can be found elsewhere\textsuperscript{12,13,14,15}. For a linear density profile, $n_e = n_{cr} (1 - x/L_n)$, the resonant fractional absorption is

$$f_A^R = 1.66 \left( \frac{\omega_p L_n}{c} \right)^{2/3} \sin^2 \phi \exp \left[ -\frac{4}{3} \frac{\omega_p L_n}{c} \sin^3 \phi \right]. \quad (\text{III.1.8})$$

Figure III.4 shows a plot of the fractional resonant absorption $f_A^R$ at different electron density scalelengths $L_n$, interacting with a 1054 nm laser beam, as a function of the angle of incidence. It can be seen that resonant absorption is predominant mostly in shorter density scalelength plasmas ($L_n < \lambda_L$). For a given density scalelength $L_n$, the laser beam angle of incidence at which it reaches its maximum value is

$$\sin \phi_{\text{max}} = 0.68 \left( \frac{\lambda_L}{L_n} \right)^{1/3}. \quad (\text{III.1.9})$$

![Graph showing resonant fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a plasma.](image)

Fig. III.4 - Resonant fractional absorption at different density scalelengths for a 1054 nm laser beam interacting with a plasma.
Fig. III.5a,b - Resonant combined with collisional scattered (non-absorbed) fraction at different electron density scalelengths for a P-polarized 1054 nm laser beam interacting with a Silicon plasma. Resonance absorption is predominant with increased target angles of incidence and shot density scalelengths (lower right region of the plot). Elsewhere, collisional absorption prevails.
Fig. III.5c - Resonant combined with collisional scattered (non-absorbed) fraction at different electron density scalelengths for a P-polarized 1054 nm laser beam interacting with a Silicon plasma. White regions have maximum transmittivity ($f_T = 1$). Black regions have no transmittivity ($f_T = 0$). Each shade of gray contour represents a 10% change on the value of the total transmittivity. Resonant absorption dominates in shorter density scalelengths and higher target angles of incidence. Collisional absorption dominates in longer density scalelengths.

In practice, collisional and resonant absorption processes operate together, and for certain plasma electron density scalelengths their combined effect cannot be discriminated. The laser light scattered fraction with both collisional and resonant absorption processes operating can be approximated as

$$f_T^{C+R} = (1 - f_A^C)(1 - f_A^R),$$  \hspace{1cm} (III.1.10)

where $f_A^R$ is applicable for P-polarization.

This product is plotted in fig. III.5a-c, for a 1 keV Silicon plasma. It is clear that the complicated shape of the scattered fraction curves and the position of their maxima (or absorption minima) are not only strongly dependent on the angle of incidence $\phi$ of the laser beam, but also in the size of the plasma electron density scalelength $L_n$. Resonant
absorption is important during the formation of short (fraction of a laser wavelength), hot plasmas ($T_e \geq 1$ keV). Collisional effects start to prevail after plasma expansion to cooler ($T_e = 100$ eV), longer density scalelengths ($L_n > 10 \lambda_L$).

Measuring the collisional plus resonant scattered fraction of a plasma with an integrating sphere set-up at different angles of incidence allows, through the use of eqns. III.1.7-10, a determination the plasma density scalelength $L_n$ and electron temperature $T_e$. A detailed description of this measurement and its results at different delay between pulses $\Delta \tau$ is in Sec. III.7.D.

B- Free Expansion of the Preformed Isothermal Plasma

Previous theoretical\textsuperscript{16,17} and experimental\textsuperscript{4,18,19} work indicates that the expansion of a plasma formed with an ultrashort laser pulse follows an isothermal model. The assumptions are:

(1)- \textit{The electrons maintain an isothermal distribution.}

The thermal conduction in the plasma is sufficiently rapid to suppress large temperature gradients. Also the heat conduction from the hot dense plasma above the critical surface can be significant in the picosecond scale, with the overdense portion of the plasma acting as a heat reservoir for the underdense portion.\textsuperscript{4,20}

(2)- \textit{The plasma has two fluid component made of warm ions and warmer electrons.}

(3)- \textit{The plasma is electrically neutral.}

The scale length of the plasma is longer than the Debye length, so charge balance between ions and electrons is maintained.

(4)- \textit{The electrons are in Boltzmann equilibrium.}

The electron response time to the electric field fluctuations, $1/\omega_{pe}$, is much shorter than the time scale of the plasma expansion, so the electron fluid is always in equilibrium with the electrostatic potential.\textsuperscript{21}

(5)- \textit{Collisional effects in the plasma are negligible.}

(6)- A planar one-dimensional expansion is assumed for plasmas of short scalelength (compared with the laser focal spot size).
With these assumptions the expansion can be described by writing the mass continuity and the momentum continuity equation for the ions and the Boltzmann equation for the electrons

\[
\frac{\partial}{\partial t} n_i + \frac{\partial}{\partial x} (n_i v_i) = 0, \tag{III.1.11a}
\]

\[
m_i n_i \left( \frac{\partial}{\partial t} v_i + v_i \frac{\partial}{\partial x} v_i \right) = -\frac{\partial}{\partial x} P_i - Z_i e \frac{\partial}{\partial x} \phi. \tag{III.1.11b}
\]

\[
n_e = n_{e0} e^{-\phi/T_e}. \tag{III.1.11c}
\]

Here \( n_i \) is the ion number density, \( v_i \) is the ion expansion velocity along \( x \), \( P_i = n_i T_i \) is the ion pressure for the ion fluid at temperature \( T_i \), \( Z_i \) is the average ion charge state, \( \phi \) is the electrostatic potential, \( n_e \) is the electron number density, \( n_{e0} \) is the electron number density at the solid target surface and \( T_e \) is the electron temperature.

**Quasi-neutrality gives**

\[
n_e = n_i Z_i. \tag{III.1.12}
\]

For an isothermal plasma, the equation for the ion pressure is

\[
\frac{\partial}{\partial x} P_i = T_i \frac{\partial}{\partial x} n_i. \tag{III.1.13}
\]

Eq. III.1.11c leads to

\[
\frac{\partial}{\partial x} \phi = -\frac{T_e}{e n_{e0}} \frac{\partial}{\partial x} n_e. \tag{III.1.14}
\]

Combining Eq. III.1.11b and Eq. III.1.11c gives:

\[
\frac{\partial}{\partial t} v_i + v_i \frac{\partial}{\partial x} v_i = -\left( \frac{Z_i T_c + T_i}{m_i n_i} \right) \frac{\partial}{\partial x} n_i = -\frac{c_s^2}{n_i} \frac{\partial}{\partial x} n_i. \tag{III.1.15}
\]

Now the similarity variable \( \zeta = (x-x_o)/t \) is introduced. At position \( x_o \) the ion number density value is \( n_{i0} \). Changing variables, we can rewrite Eq. III.1.11a and Eq. III.1.15 as

\[
(v_i - \zeta) \frac{\partial}{\partial \zeta} n_i = -n_i \frac{\partial}{\partial \zeta} v_i, \tag{III.1.16a}
\]

\[
(v_i - \zeta) \frac{\partial}{\partial \zeta} v_i = -\frac{c_s^2}{n_i} \frac{\partial}{\partial \zeta} n_i. \tag{III.1.16b}
\]
Eliminating the derivatives of the ion number density and expansion velocity, the self-similar solution for the expansion velocity is

\[ v_i = \zeta \pm c_s. \]  (III.1.17)

Substituting equation III.1.7 in equation III.1.6b and integrating the result in \( \zeta \), the self-similar solution for the ion number density is

\[ n_i = n_{i0} e^{-v_i/c_s}. \]  (III.1.18)

Figure III.6 shows a plot of the ratio \( n_i/n_{i0} \) and the expansion Mach number \( M = v_i/c_s \) associated to the planar solutions equation III.1.17 and equation III.1.18.

\[ \text{Fig. III.6} \text{- Ion density ratio} \left( n_i/n_{i0} \right) \text{ and Mach expansion speed for planar isothermal expansion of a plasma.} \]

For the spherical expansion of an isothermal plasma a procedure similar to the one used to find the plane solution is followed. It leads to the following ordinary differential equations\textsuperscript{22,23,24}

\[ \left[ (M - \Xi)^2 - 1 \right] \Xi \frac{\partial M}{\partial \Xi} = 2. \]  (III.1.19a)
\[(M - \Xi) \frac{\partial M}{\partial \Xi} = -\frac{1}{n_i} \frac{\partial n_i}{\partial \Xi}, \quad (\text{III.1.19b})\]

where \(M = v/c_s\) is the expansion Mach number and \(\Xi = (r-r_0)/c_st\) is the radial similarity variable.

The self-similar solution for the ion number density is the fast decreasing Gaussian

\[n_i = n_{i0} e^{-(\Xi-1)^2}. \quad (\text{III.1.20})\]

Substituting Eq. III.1.20 in Eq. III.1.19b and using this result to eliminate \(\partial M/\partial \Xi\) in III.1.19a, the self-similar solution for the expansion velocity is

\[M = \Xi + \frac{1}{2\Xi} \pm \sqrt{\frac{1}{4\Xi^2 + 2}}. \quad (\text{III.1.21})\]

Figure III.7 shows a plot of the ratio \(n_i/n_{i0}\) and the expansion Mach number \(M = v_i/c_s\) associated to the spherical solutions Eq. III.2.10 and Eq. III.1.21.

**Fig. III.7** - Ion density ratio \((n_i/n_{i0})\) and Mach expansion speed for spherical isothermal expansion of a plasma.
The expansion of the plasma is more accurately modeled as a combination of the planar and spherical solutions (see schematic on fig III.24). The planar solution describes the plasma expansion better for distances to the target surface closer than the laser spot diameter. The spherical solution describes the expansion of the plasma farther from the target surface.

**C - Plasma Heating by the Interaction Pulse**

Even when the expansion of the preformed plasma is isothermal, its thermodynamic state can be different from the isothermal prediction because of heating during the interaction with the second pulse. Both resonant absorption heating near the critical surface and collisional absorption heating in the expanded volume of the plasma are present. This heating effect will increase the SBS threshold intensity in the local resonant region, due to the increase of the local ion-sound speed and the steepening of plasma density, velocity and temperature gradients.

The electromagnetic wave transfers energy to the electrons at a collisional damping rate

\[
\nu_{\text{coll}} = \nu_{\text{ei}} \frac{\omega_{\text{pe}}^2}{\omega_0^2} = \nu_{\text{ei}} \frac{n_e}{n_{\text{cr}}}.
\]  

(III.1.22)

As the product of the collisional damping rate \(\nu_{\text{coll}}\) with the laser pulse width \(\tau\) is typically

\[
\nu_{\text{coll}} \tau = \frac{n_e}{n_{\text{cr}}} \nu_{\text{ei}} \tau \ll 1.
\]  

(III.1.23)

the electrons keep most of the collisional absorbed electromagnetic energy. The laser pulse duration is too brief to allow the electrons to equalize their thermal energy with the ion species at the same time the EM wave is being absorbed and Brillouin scattered. Then the local variation on electron temperature \(\Delta T_e\) is much larger than that of the ion temperature \(\Delta T_i\),

\[
\Delta T_e \gg \Delta T_i.
\]  

(III.1.24)

Collisional and resonant absorption heating is greater near the critical surface than in the expanded tail of the preformed plasma, meaning greater electron temperature and higher average ion charge in this region. The local variation of electron temperature \(\Delta T_e\) is expected to have a negative gradient along the line of plasma expansion.
\[ \frac{\partial}{\partial z}(\Delta T_c) \leq 0. \quad (\text{III.1.25}) \]

Consequently, the local ion-sound speed \( c_s \) increases in this region. Therefore, the SBS threshold intensity increases near the critical surface, and the region of resonance moves to the cooler, more rapidly expanding parts of the plasma. This translates into larger Doppler blueshifts for the backscattered Brillouin spectra.

**D- LILAC Simulations**

Simulations of the plasma behavior with realistic experimental parameters were performed using the one-dimensional Lagrangian hydrodynamic code LILAC, for both planar and spherical expansion profiles.\(^{25}\) This code includes ray tracing, Thomas-Fermi equation of state and multigroup diffusion radiation transport. The ray tracing is carried out using the azimuthally averaged spatial profile of a typical beam with the inverse bremsstrahlung opacity corrected for the Langdon effect.\(^{26}\) The electron thermal energy is transported using a flux-limited diffusion model in which the effective flux was defined as the minimum of the diffusion flux and the free-streaming flux (the sharp cutoff method). The opacities used in the radiation transport calculation are obtained from the Los Alamos LTE (Local-Thermodynamic-Equilibrium) astrophysical library.\(^{27}\)

Figures 8a-d show the results of a typical LILAC simulation for the planar expansion of a Silicon plasma at times just before and after of the interaction beam appearance. This plasma has been preformed by a laser pulse of intensity \( I_{\text{pre}}=10^{15} \, \text{W/cm}^2 \), followed 260 ps later by an interaction pulse of intensity \( I_{\text{int}}=10^{16} \, \text{W/cm}^2 \). The initial plasma spot at the laser focus is 16 \( \mu \text{m} \times 16 \, \mu\text{m} \), for a target radius of curvature of 1000 \( \mu\text{m} \). The solid lines show the results when no plasma heating by the interaction pulse has been included. The dashed lines plot the results of plasma heating by the interaction pulse; the time at the legend indicates how long the interaction pulse has been heating the plasma. As the interaction laser pulse has a FWHM of 1.6 ps, it takes it at least 3.2 ps to traverse its path in and out of the subcritical regions of the plasma. The ratio between electron number density \( n_e/n_{\text{cr}} \) over the critical density is plotted in fig. 8a. It can be seen that interaction pulse heating only changes the density profile in a small region close to the critical surface. The profile of the plasma expansion speed \( v_{\text{exp}} \) is shown in fig. 8b. The expansion speed changes slightly with interaction pulse heating, the change greater close to the plasma critical region.
**Fig. III.8a - LILAC calculated \( n_e/n_{cr} \) profile (planar).**

**Fig. III.8b - LILAC calculated plasma (planar) expansion speed profile.**

The Silicon plasma is preformed by a laser pulse of intensity \( I_{lo}=10^{15} \text{ W/cm}^2 \), followed 260 ps later by an interaction laser pulse of intensity \( I_{int}=10^{16} \text{ W/cm}^2 \). A Solid line indicates no plasma heating by the interaction pulse. Dashed lines indicate plasma heating by the interaction pulse. The time at the legend indicates how long the interaction pulse has been inside the plasma; a 1.6 ps FWHM laser pulse takes at least 3.2 ps to enter and leave completely the plasma.
Fig. III.8c - LILAC calculated electron temperature profile (planar).

Fig. III.8d - LILAC calculated ion temperature profile (planar).

The Silicon plasma is preformed by a laser pulse of intensity $I_{\text{pre}} = 10^{15}$ W/cm$^2$, followed 260 ps later by an interaction laser pulse of intensity $I_{\text{int}} = 10^{16}$ W/cm$^2$. A Solid line indicates no plasma heating by the interaction pulse. Dashed lines indicate plasma heating by the interaction pulse. The time at the legend indicates how long the interaction pulse has been inside the plasma; a 1.6 ps FWHM laser pulse takes at least 3.2 ps to enter and leave completely the plasma.
The longer the time after the appearance of the interaction pulse inside the plasma, the larger is the region where $v_{\text{exp}}$ is increased and the amount of the change. The electron temperature is plotted in fig. 8c. The electron temperature increases dramatically when interaction pulse effects are considered. The increase covers the entire length of the expanded plasma—but is greater close to the denser critical region, and it grows in time after the interaction pulse. The ion temperature, plotted in fig. 8d, only increases slightly when heating effects are considered. This behavior clearly shows that electron-ion collisionality effects start to be noticeable only after relatively long periods of time (compared with pulse duration) in this plasma regime ($v_{\text{ci}} \tau \ll 1$).

LILAC simulations for planar and spherical expansion were done also for plasmas preformed by a laser pulse of intensity $I_{\text{se}}=10^{15}$ W/cm$^2$, but with delay between pulses of 260, 510, 760, 1010 and 1260 ps for Silicon and 1010 ps for Gold and Parylene-N. In all cases the interaction pulse intensity was $I_{\text{im}}=10^{16}$ W/cm$^2$. The overall behavior of relevant parameters of these plasmas is similar to the 260 ps Silicon case plotted in fig. 8a-d.

Figure 9a-e shows comparative plots for electron number density ratio, plasma expansion velocity, electron and ion temperature, and ionization number for the planar expansion of Silicon, with delays between preforming and interaction pulses of 260, 510, 760, 1010 and 1260 ps. The plots show the change experimented by the spatial profiles of these parameters due to the effects of plasma expansion after formation by the first laser pulse. In all cases plotted, the subsequent interaction pulse heating has been taken in account.
Fig. III.9a - LILAC calculated $n_e/n_{cr}$ profile (planar).

Fig. III.9b - LILAC calculated plasma expansion velocity profile. The Silicon plasma is preformed by a laser pulse of intensity $I_{0,l}=10^{15}$ W/cm$^2$, followed by an interval $\Delta t$ later by an interaction laser pulse of intensity $I_{int}=10^{16}$ W/cm$^2$. 
Fig. III.9c - LILAC calculated electron temperature profile.

Fig. III.9d - LILAC calculated ion temperature profile.

The Silicon plasma is preformed by a laser pulse of intensity $I_{\text{p}}=10^{15}$ W/cm$^2$, followed by an interval $\Delta \tau$ later by an interaction laser pulse of intensity $I_{\text{in}}=10^{16}$ W/cm$^2$. 
III.2 - TTT Chirped Pulse Amplification Laser System

The backscattered light measurements were performed with the TableTop Terawatt (TTT) Chirped Pulse Amplification laser system, at LLE. A detailed description of this system can be found in Refs. 28, 29. Figure III.10 shows a diagram of the TTT laser system.

The pulse train from the Nd:YLF oscillator is made up of a 100 MHz, 50 ps mode-locked pulses at a wavelength of 1053 nm and 0.3 Å bandwidth. These pulses pass through a 800 m single mode optical fiber, which broadens the bandwidth to approximately 40 Å, due to the self-phase modulation and imposes a chirp in the pulse, with bandwidth of approximately 40 Å. Expansion gratings then are used to stretch the pulse to 500 ps. One of the train of pulses is then amplified to approximately 0.5 mJ in a regenerative amplifier. Using Pockels cells and fast electronics, a single gain-narrowed 300 ps pulse is switched out of the train, spatially filtered and sent to a 9 mm diameter rod amplifier, where its energy is raised to ~50 mJ. A 3:1 spatial filter then removes any noise or intensity spikes which may appear and the cleaned pulse is sent through a 30 mm
diameter single-pass rod amplifier, reaching energies up to 1.5 J. The pulse is spatially filtered (1.2:1) again, expanded through a 8/3:1 Galilean beam expander and recompressed with another grating pair to approximately 1.6 ps, with energies up to 1.5 J.

Fig. III.10 - The TTT Chirped Pulse Amplification laser system.

III.3 - Pulse Splitting, Delaying and Stacking

As described in Chapter II, for convective SBS to be above threshold for picosecond pulses, the plasma density scalelength must be longer than the laser wavelength. The plasma must expand for much longer than the 1.6 ps pulse duration for typical ion-sound
speeds in the $10^6$-$10^7$ cm/s range. In this experiment, SBS is studied for density scalelengths of 10 to 100 laser wavelengths, which require expansion times between 200 to 1500 ps.

Suitable plasmas can be created using two laser pulses separated by 200 ps or more. The first one is the plasma preforming pulse followed after a delay $\Delta \tau$ by the higher intensity interaction pulse, which will scatter in the preformed plasma. The magnitude of $\Delta \tau$ is one of the three most important parameters in the experiment—the other two are the intensities of the preforming and interaction pulses at the target surface. The state of the plasma at the interaction time determines the SBS threshold.

Figure III.11 shows a diagram of the optical set up used to split the incoming single pulse laser beam in two pulses, to delay and to set the intensity ratio of these pulses one with respect to the other and to stack them in a double pulse train laser beam. The set-up is placed after the expansion gratings, just before the beam enters the regenerative amplifier, to ensure that the spatial modes and pointing of the two pulses are identical. A Michelson interferometer set-up is used to introduce a delay between pulses proportional to the interferometer’s arms length difference. Also the contrast ratio between the plasma preforming and interaction pulses is set through the manipulation of their polarization components.

The S-polarized ($90^\circ$ with respect to the horizontal) beam coming from the oscillator/compression gratings is injected through a 50% glass beam splitter. Half of the light goes straight along the shortest arm—of length $L_1$. This will become the plasma preforming pulse. The pulse passes through a $\lambda/2$ waveplate, is rotated an angle $\phi_1$ with respect to the vertical, is reflected back with a fixed mirror, and passes again through the $\lambda/2$ waveplate. Now the pulse has a mixed polarization—$(2 \cos^2 \phi_1 - 1)$ S-component and $\sin 2\phi_1$ P-component. It is reflected again by the beam splitter out of the interferometer. The P-component of this pulse is rejected by a cube polarizer before entering the first Pockels cell. After that, it is injected into the regenerative amplifier.

The second half of the original S-polarization beam is reflected by the splitter $90^\circ$ along the longer arm—of variable length $L_2 \geq L_1$. This will become the interaction pulse. The pulse passes through a second $\lambda/2$ waveplate, is rotated an angle $\phi_2$ with respect to the vertical, is reflected back by a mirror mounted on a movable translation stage, and passes again through the second $\lambda/2$ waveplate. After transmission through the splitter, it is injected into the regenerative amplifier, delayed behind the first (preforming) pulse an
interval $\Delta \tau = 2(L_2 - L_1)/c$. The P-component of the delayed (interaction) pulse also is rejected by a cube polarizer before entering the first Pockels cell.

![Diagram of pulse splitting, delaying, and stacking set-up.](image)

**Fig. III.11 - Pulse splitting, delaying and stacking set-up.**
The S-polarized pulse train from the Nd:YLF oscillator comes from the left of the diagram. The vertical arm of the Michelson set-up generates the plasma preforming pulse; the horizontal arm generates the delayed interaction pulse. The S-polarized double pulses exit to the right into the regenerative amplifier.

The two S-polarized pulses separated by an interval $\Delta \tau$ ($0 \leq \Delta \tau \leq 1500$ ps) and different amplitude contrast ratio ($0.1 \leq I_{int}/I_{for} \leq 10$) leave the splitter to be amplified and compressed.

### III.4 - High-Contrast Laser Pulses

After leaving the compression grating, a laser pulse is a Gaussian superimposed on a $\sim 100$ ps duration pedestal. This pedestal is generated by the unchirped parts of the laser pulse emanating from the fiber. The pedestal has a considerable amount of energy (approximately $1/10$ of the maximum main peak energy). In a single pulse it can preform a small plasma before the interaction of the main (Gaussian-like) portion of the pulse.
To eliminate this pedestal a saturable absorber cell is inserted in the laser beam path after it leaves the double-pass compression gratings of the TTT CPA laser (see lower part of fig. III.10). An aniline (Kodak Q-switching dye #9860) dissolved in nitrobenzene (~5·10⁻⁵ M), acts as a saturable absorber. For low laser pulse energies the saturable absorber has a low light transmittance (less than 0.2% below ~10¹² W/cm²). For higher intensities, the solution light absorption capability saturates and allows the transmission of ~35 to 40% of the incoming light. A 1-cm thick glass cell on a kinematic mount contains the solution. Figure III.12 shows the measured transmittance of the saturable absorber cell and its corresponding semi-empirically fit curve.³²

An autocorrelator consisting of a LiIO₃ crystal and a linear photodiode array (Reticon RC105) measures the output pulse duration. Figure III.13 shows an autocorrelation trace³³,³⁴ of a typical low-contrast pulse generated by the laser system (hairline dashed line). The second trace in fig. III.13 shows an autocorrelation trace of a typical high-contrast pulse (hairline full line) produced after the passage of a low-contrast pulse through a saturable absorber.

![Graph showing saturable absorber transmission curve for the laser beam.](image)

**Fig. III.12 - Saturable absorber transmission curve for the laser beam.**

After the passage through the saturable absorber cell, the laser pulses are Gaussian like, approximately 1.2 to 1.6 ps in length, with the pedestal intensity suppressed to values less
than $10^{-5}$ times the main peak intensity (see fig. III.13). Thus a preformed plasma under uncontrollable conditions no longer appears.

**Fig. III.13 - TTT laser pulse before (low contrast) and after (high contrast) the passage through the saturable absorber.**

The two small peaks marked as Etalon effects are instrumental artifacts, coming from optical elements with parallel surfaces in the autocorrelator used in the measurement of the laser pulse temporal profile.

### III.5 - THE ZETA-TANK

The experimental SBS measurements were made using the Zeta vacuum chamber. Figure III.14 shows a diagram of the experiment's light path geometry and the placement of the diagnostics used to measure the parameters of both incident and backscattered light.
Fig. III.14 - Diagnostics set-up to measure both incident and backscattered light at TTT Zeta tank.
The solid line shows the path of the incoming double-pulse laser beam into the Z-tank. The dashed line shows the path of the SBS backscattered light out to the diagnostics.

A 60 cm focal length positive lens focuses the beam into the center of the Zeta tank. Figure III.15 shows a cross-section of the focal intensity profile, taken in the plane normal to the direction of beam propagation. The lens focal spot normal to the propagation of the beam is an ellipse (because of uncorrectable astigmatism from the saturable absorber cell)
of semiaxis $11\pm0.6 \, \mu m \times 17\pm1 \, \mu m$ ($1/e$ intensity) for the 2.4 cm radius incoming beam with a $f_\# = 12$. The corresponding Rayleigh range is approximately $300\pm30 \, \mu m$, nine times longer than the laser spot size, ensuring that nearly planar wavefronts irradiate the target. The Rayleigh range is longer than both the density and the velocity scalelengths of the plasma in the resonance region. Polarization mixing and uncertainty in the angle of incidence are avoided.

![Image of laser focus](image)

**Fig. III.15 - Normal transverse image of the laser focus (far field).**

The enclosed regions are isointensity contours of the focused laser spot far field image (in the plane normal to the direction of beam propagation), taken inside the Zeta vacuum chamber. Each shade of gray contour represents a $\sim 7\%$ change on the beam's intensity. This beam has passed through the saturable absorber cell. The image size is 60 pixels $\times$ 60 pixels and the scale is 1.84 pixel/$\mu m$.

A $\lambda/2$ waveplate before the lens allows the beam linear polarization to be continuously changed between $S$- and $P$-. A coaxiled He-Ne (green) laser beam is used for the alignment of targets and diagnostics.

A motorized, computer-controlled three-axis linear translation stage holds the targets at the best focus of the Zeta lens. The translation of the target can be done manually or be computer-controlled. The target support is electrically insulated from ground. A CCD camera/monitor placed normal to the laser propagation direction provides visual feedback to manually control the position and orientation of the target, and to examine the target surface conditions during the experiment.
Most of the measurements are done at 55° laser angle of incidence (ϕ) with respect to the target normal. This angle has the maximum resonant absorption for \( p \)-polarized plasma formation laser pulse (when the density scalelength is \( L_n < 1 \mu m \))\textsuperscript{36}. In the density scalelength measurements, the target was rotated to angles between 22.5° to 82.5° with respect to the direction of the incoming beam. A target rotator incorporating a computer controlled small stepper-motor allows the target to be rotated clockwise and anticlockwise in 7.5° steps with minimal error. For experiments done with the target surface at angles of incidence \( ϕ \) greater than zero (22.5° ≤ \( ϕ \) ≤ 82.5°), the effective target area is a larger ellipse of semiaxis \( (11/\cos \phi)\pm 0.6 \mu m \times 17\pm 1 \mu m \) (1/e intensity).

The target chamber base pressure is kept below 10\(^{-6}\) Torr to avoid the formation of an air plasma in front of the target and to prevent the charge exchange between plasma ions and the ambient gas atoms.

### III.6 - Targets

The targets are solid planar strips approximately one inch long and a 1/4 inch wide. The size of the target accommodates ~120 shots, with each laser shot crater separated by ~1 mm (twice the crater size). All SBS related laser shots interacted with a undamaged fresh surface. Three types of solid targets were used:

- **Polished silicon** (\( Si, A = 28.1, Z = 14 \))
  
  One millimeter thick pure amorphous polished silicon planar strips were the primary targets used for SBS measurements. Silicon was chosen due to its moderately low Z-number, availability and high optical quality of the target surface.

- **Gold** (\( Au, A = 197, Z = 79 \))
  
  A 2000 Å thick gold film was deposited over a flat glass or silicon substrate. Gold was selected to measure the effects of high mass and ion charge in SBS. It also provides an excellent quality optical surface.

- **Parylene-N** (\( CH, A = 12+1, Z = 6+1 \))
  
  A 5000 Å thick Parylene-N film, was deposited over a flat glass or silicon substrate. Parylene-N was selected to measure the effects of a multispecies plasma on SBS. The mix of a very low Z-number element (H) with a moderately low Z-number element (C) produces a plasma with multiple ion-acoustic modes and enhanced Landau damping\textsuperscript{37}. Great care must be taken to maintain the optical conditions of the surface, as the
Parylene-N film has a tendency to shrink, be scratched and break, exposing the glass substrate.

For both Gold and Parylene-N, the ablation depth of the laser crater (~1000 Å) is much less than the thickness of deposited film.\textsuperscript{38}

\section*{III.7 - Diagnostics}

The total energy and relative amplitude of the incident laser pulses are measured simultaneously before entering into the Z-tank. The total and spectral energy of the backscattered light plus the spectrum of the backscattered light are measured. Figure III.14 shows a diagram of the experiment’s light path geometry and the placement of the diagnostics used to measure the parameters of both incident and backscattered light.

\subsection*{A- Incident Laser Light}

\subsubsection*{a- Total incident laser energy}

The total incident energy of the incoming laser beam is measured after the saturable absorber cell, using a large area UDT-100 PIN diode, biased at +9 V. The diode is placed behind the second turning mirror—with a transmission of approximately 2.5\%, and a condenser lens. This diode is filtered with a RG-1000 short wavelength cut-off filter for the 1053 nm laser wavelength. Additional filtering is used to limit the size of the signal in the PIN diode, protecting it from damage and saturation.

\subsubsection*{b- Contrast between preforming and interaction pulses}

The contrast between the preforming and interaction pulses is measured by an ultrafast diode (rise time $\approx 75$ ps), coupled via an optical fiber to a fast transient scope (Tektronix SCD-5000).\textsuperscript{39} The diode is placed behind the first turning mirror, with transmission of approximately 2.5\%, with no condenser lens, so that it samples only a small part of the beam. This set-up allows pulses separated by between 75 and 1500 ps and contrast ratios between 1 and 15 to be measured. Figure III.16 shows fast transient scope traces of a single pulse and a double pulse train separated by a delay of 530 ps.
Fig. III.16a - Transient-Scope trace for a single laser pulse.

Fig. III.16b - Transient-Scope trace for a double laser pulse, separated by 530 ps.

Pulse contrast $I_{int}/I_{for}$ was measured by taking the ratio between the second over the first voltage, after checking that the separation between their peaks was set at the proper delay.
B- Backscattered Light

The backscattered light originates close to the focus of the Z-tank focusing lens, in the center of the vacuum chamber. Restricted by the Z-tank window diameter to a collection solid angle of 0.05π steradians, this light is collected in a collimated beam by the Z-lens, being relayed back along the same optical path as that of the incident laser beam (see the dotted line in Figure III.14).

A coated glass flat, of reflectivity 1.8 %, located after the output of the compression gratings collects part of the backscattered energy coming from the Zeta tank and diverts it to a antireflection coated 200 cm focal length positive lens. This lens collects and focuses the backscattered light to be analyzed spectrally.

a- Total backscattered energy

The total energy of the backscattered light is measured with a large area UDT-100 PIN diode, biased at +9 V, placed at the focus of the collection lens. This diode is filtered with a RG-1000 pass-band filter for the 1053 nm laser wavelength. Additional filtering is used to limit the size of the signal in the PIN diode, protecting it from damage and signal saturation.

b- Time-integrated backscattered spectrum

The backscattered light spectrum integrated over the interaction time of the pulse with the plasma is measured with a 1 m Czerny-Turner spectrometer (using a Bausch & Lomb flat grating with 1200 groves/mm, blazed at 17°).

The collected backscattered light is injected through a short focal length condenser lens and a glass diffuser into the spectrometer. The resolved spectra is detected by an EG&G/PARC Model 1471A Optical Multichannel Analyzer/Photo Diode (OMA/PD) and converted and stored digitally with a GPIB/PC set-up.

The calibration of the spectrometer was done by taking the known spectrum of a Rubidium lamp. Figure III.17 shows a plot of the Nd:YLF oscillator line spectra taken at different spectrometer entrance slit apertures. The input spectral line width is approximately 0.5 Å. The different results at different slit apertures allows the measurement of the spectrometer resolution. The integration time settings for the OMA/PD were the same used in the experiment (t_{OMA} = 30 ms). The best resolution—1.2 Å—was for a slit aperture of
0.14 mm (which agrees very well with the theoretical value for the slit size to achieve maximum resolution,\textsuperscript{41}

\[ \text{slit} = \sqrt{2\lambda_L f_{CT}}, \]

with the laser wavelength \( \lambda_L = 1053 \text{ nm} \) and the Czerny-Turner spectrometer focal length \( f_{CT}= 1 \text{ m} \). For narrower slits the signal amplitude decreased without changes in the measured spectral width.

![Graph showing different slit apertures](image)

**Fig. III.17 - Nd:YLF Oscillator spectra for different slit apertures.**
The Nd:YLF oscillator input (calibration) signal has a spectral width of 0.5 Å.

**C- Energy calibration of the Z-tank diagnostics**

The total incoming laser energy PIN diode, the total backscattered energy PIN diode and the spectral energy measured by integration of the digitized OMA/PD spectra are calibrated in one operation. The diodes give their energy measurements in counts, and the integrated signal in the spectrometer appears in counts-nm. The calibration of these devices allows conversion of the instruments’ units of energy measurement to energies expressed in milliJoules (mJ).
Figure III.18 shows a schematic of the calibration set-up used. The calibration device is an absolutely calibrated Joule-meter (Molelectron G-50), placed inside the Zeta tank, close to the window (to reduce damage to the device surface and filters by the high intensity incident laser beam if placed near the focus). A 50±1% reflectivity$^{42}$ mirror is set normal to the path of the incident laser beam. In this way, the same amount of energy arrives to the Total Incident Energy PIN diode/Molelectron Joule-meter transmitted path and to the Total Backscattered Energy PIN diode/Spectrometer reflected path. The laser is fired at different energy levels and calibration formulas are found relating the diagnostic instruments measurement values to the Joule-meter measured energies. These are

$$E[\text{mJ}]=(13.8 \pm 0.1) \times 10^{-3} \cdot \text{BPD[counts]}; \ R = 0.9953$$

$$E[\text{mJ}]=(7.57 \pm 0.05) \times 10^{-3} \cdot \text{OMA[counts-nm]}; \ R = 0.9947$$

**Fig. III.18 - Calibration set-up for all energy diagnostics at the Z-tank.**

The continuous line shows the path of the incident double-pulse laser beam into the Z-tank. The dashed line shows the path of the light reflected to the diagnostics. The units in square brackets at the side of the calibrated parameters indicate the instrument output measurement units. All energy parameters are referred in millijoules after calibration.
Fig. III.19 a, b- Calibration plot for total and spectral backscattered energy. The upper plot (a) shows the calibration of the total (time-integrated) backscattered energy measured using the large area Back PIN diode. The lower plot (b) shows the calibration of the spectral (time-integrated) energy measured using the OMA/PD and spectrometer. The formula over the plot gives the conversion from the instrument units [counts for the BPD, counts-nm for the OMA] to physical units [milliJoules]. The error bars for the measurements in both axis (10 counts for the BPD, 50 counts-nm for the OMA, 1 mJ for the Joulemeter) are comparable to the size of the plot point symbol for both plots.
Figure III.19 shows the calibration curves for the Total (time-integrated) Backscattered Energy PIN diode (plot a), the spectral energy (time-integrated) measured by integration of the spectral area obtained by the OMA/PD (plot b), and a plot of the PIN data plotted versus the OMA data (plot c), taken simultaneously for single laser pulses, with the corresponding conversion formulae from instrumental to physical units. The plot's horizontal axis always shows the instrument (Back PIN or OMA/PD) units, while the vertical axis shows the energy units, because the Molecron joulemeter is used only for calibration. In an experimental run, the instruments' units must be converted to energy units, not the reverse.

![ Calibration plot for total and spectral backscattered energy. ]

The plot shows the linear correlation between the calibrations of the total (time-integrated) backscattered energy measured using the large area Back PIN diode the spectral (time-integrated) energy measured using the OMA/PD and spectrometer. The error bars for the measurements in both axis (10 counts for the BPD, 50 counts-nm for the OMA) are comparable to the size of the plot point symbol for both plots.

D- Density scalelength

The electron density scalelength of the plasma \( L_n \) is estimated by measuring the collisional plus resonant absorption/reflectance of the laser light at different combinations of P- and S-polarizations at different incident angles with an integrating sphere surrounding the target in the center of the Zeta tank.\(^{43} \) This is a hollow plastic sphere,\(^{44,45} \) covered in its inner surface with an approximately 1 mm thick coat of diffusing Kodak 6080 white reflectance paint.\(^{46} \) This device has nearly \( 4\pi \) collection angle and the measured scattered
intensity is independent of the angular radiation distribution of the source. Four large area UDT-100 PIN diodes (biased at +9 V) with different filter values (including one short-wavelength cutoff RG-1000 for 1054 nm) are located on the surface of the sphere to measure the scattered light.

Figure III.20 shows a diagram of the integration sphere set-up used to measure density scalelengths. The plasma formation prepulse is incident on the solid target, creating a near-solid density plasma of very short electron density scalelength ($L_{n,for} < \lambda_L$). After an interval $\Delta \tau$, the preformed plasma expands to a longer electron density scalelength ($L_{n,int} < \lambda_L$), then the interaction pulse arrives. Figure III.1 shows a schematic of the double-pulse process of formation and expansion of a plasma. The fraction of the total (formation plus interaction pulses) non-absorbed (reflected plus scattered) light is measured. The target is rotated with a different angle $\phi$ ($22.5^\circ \leq \phi \leq 82.5^\circ$, measured between the target normal and the laser beam line of propagation) for each laser shot.

![Diagram of integration sphere set-up](image)

**Fig. III.20 - Integrating Sphere set-up to measure density scalelengths $L_n$ by maximum resonant absorption.**

The sphere is mounted surrounding the target in the center of the Zeta tank.

The non-absorbed fraction for the formation (of intensity $I_{for}$) plus the interaction (of intensity $I_{int}$) pulses at a target angle $\phi$ can be calculated, for a pulse contrast ratio $CR = I_{int}/I_{for}$, delay between pulses $\Delta \tau$, and polarization values, as
\[ f_{T,\text{tot}}^{C+R}(\phi; \Delta \tau, \text{pol}) = \frac{1}{1 + CR} \left[ f_{T,\text{for}}^{C+R}(\phi; 0, \text{pol}) + CR f_{T,\text{int}}^{C+R}(\phi; \Delta \tau, \text{pol}) \right]. \quad (\text{III.7.1}) \]

In § III.1, it was shown that the non-absorbed (scattered) fraction \( f_T^{C+R} \) of a single pulse (P-polarized) of laser light with both collisional and resonant absorption processes operating is:

\[ f_T^{C+R} = (1 - f_A^C)(1 - f_A^R), \quad (\text{III.1.10}) \]

with

\[ f_A^C = 1 - \exp \left( -\frac{32 \, \nu_{eL_n}}{15 \, c} \cos 5 \theta \right), \quad (\text{III.1.5}) \]

and

\[ f_A^R = 1.66 \left( \frac{\omega_0 L_n}{c} \right)^{2/3} \sin^2 \theta \exp \left[ -\frac{4 \, \omega_0 L_n}{3 \, c} \sin^3 \theta \right] \quad (\text{III.1.8}) \]

This single pulse (P-polarized) scattered fraction \( f_T^{C+R} \) is plotted in fig. II.5a-c for an 1keV Silicon plasma at different scalelengths and angles of incidence.

Fig III.21 shows the integrating sphere (plus backscattered light exiting through the entrance window, measured with the PIN diode) measured scattered fraction of a Silicon plasma for a single, P-polarized laser pulse at different target angles. Fitting the measured data points to the scattered fraction formula III.1.10, the values of the plasma density scalelength \( L_n \) and electron temperature \( T_e \) (with their errors) can be obtained correspondingly from the fit dimensionless parameters \( \omega_0 L_n/c \) and \( \nu_{d*}/\omega_0 \).

The same non-absorbed light measurement is shown in Fig III.22, but now for double laser pulses separated by 1260 ps. In this case, the measured data points were fitted to the scattered fraction formula III.7.1. The contrast ratio between interaction and preforming pulses was 7. Also, the values of \( (\omega_0 L_n/c)_{\text{for}} \) and \( (\nu_{d*}/\omega_0)_{\text{for}} \) for the plasma preformation pulse are taken from the single pulse values found in the fit (using eqn. III.1.10, fig. III.21). The values of the plasma density scalelength \( L_n \) and electron temperature \( T_e \) (with their errors), at the time when the interaction pulse is present in the plasma, can again be obtained from the fit parameters \( (\omega_0 L_n/c)_{\text{int}} \) and \( (\nu_{d*}/\omega_0)_{\text{int}} \).
Fig. III.21 - Integrating Sphere measured scattered fraction for a single laser pulse ($\Delta t = 0$ ps) at different Silicon target angles. The scattered fraction fit curve corresponds to formula III.1.10.

Fig. III.22 - Integrating Sphere measured scattered fraction for double laser pulses ($\Delta t = 1260$ ps) for different Silicon target angles. The scattered fraction fit curve corresponds to formula III.7.1. The contrast ratio between interaction and preforming laser pulses is CR= 7.
Figure III.23 shows the integrating sphere electron density scalelength measurements for a Silicon plasma interacting with double laser pulses, separated by delays $\Delta \tau$ between 510 and 1260 ps, compared to the LILAC calculated values (the error bars in LILAC data points appear because these values were found by fitting eqn. III.18 to planar expansion simulations, and eqn. III.1.20 to spherical expansion simulations). All of them were collected proceeding similarly to the electron density scalelength value for $\Delta \tau = 1260$ ps found in fig. III.22; the corresponding values of $(\omega_0 L_n/c)_{\text{int}}$ and $(v_\alpha^*/\omega_0)_{\text{int}}$ were found by non-linear curve-fitting of eqn. III.7.1, for pulse contrast ratio $CR = I_{\text{int}}/I_{\text{for}} = 7$, and using the values of $(\omega_0 L_n/c)_{\text{for}}$ and $(v_\alpha^*/\omega_0)_{\text{for}}$ as those found for the single pulse for the first term in III.7.1.

![Graph showing Silicon density scalelengths](image)

**Fig. III.23 - Measured and LILAC-calculated electron density scalelengths $L_n$ at different delays $\Delta \tau$ for Silicon planar targets.** The error bars in LILAC data points appear because these values were found by fitting eqn. III.18 to planar expansion simulations, and eqn. III.1.20 to spherical expansion simulations.

The integration sphere density scalelength measurements agree closely with the planar expansion LILAC simulations, $L_n \propto c_s \Delta \tau$, described by the solutions III.1.17-18 to the system III.1.11. The region of resonantly absorbing plasma is then very near to the critical
surface of the planar silicon targets. The expanding plasma flow is normal when near to the flat target original surface and to the also planar plasma’s critical surface, at least by distances comparable to the size of the laser focus (~12 $\mu$m $= 12 \lambda_L$, see § III.5 and fig. III.15); at farther distance from the original target surface the plasma expands spherically, propagating outwards radially from the center of the laser focus. A schematic of the appearance of the plasma expansion is shown in fig III.24.

The electron temperatures near the plasma critical surface, derived from these scalelengths and the $v_\phi/\omega_0$ values measured with the integrating sphere, are shown at fig. III.25. The I.S. measured values are at least one order of magnitude higher than those from the corresponding LILAC planar expansion simulations without heating, but a better correlation for the LILAC planar expansion simulations that consider the effect of plasma heating by the interaction pulse. As the I.S. measured temperatures are derived from parameters ($v_\phi/\omega_0$) associated to nonlinear curve fits of collisional plus resonant absorption’s scattered fraction curves (eqn III.1.10 for single-pulse, eqn. III.7.1 for double-pulses) like in figs III.21-22, a great uncertainty in their final magnitude results.

Table III.1 compiles the diverse scalelengths and values of the relevant plasma parameters for both LILAC calculations and integration sphere measurements for Silicon, at different interval between preforming and interaction pulses $\Delta t$.

---

**Fig. III.24 - The expansion of a plasma near a solid target surface.**

The plasma expands normally to the target surface for distances less than the laser focus spot size. The plasma critical surface and the region of maximum resonant absorption are inside this planar expansion region. Farther away it expands spherically.
Fig. III.25 - Electron temperature $T_e$, Integration sphere measured near the critical surface and LILAC-calculated average, at different delays $\Delta \tau$ for Silicon planar targets.

LILAC data is for planar expansion simulations.

Table III.1 - LILAC calculated plus Integration Sphere measured Scalelengths near the plasma critical surface.

<table>
<thead>
<tr>
<th>TARGET</th>
<th>$\Delta \tau$ [ps]</th>
<th>$L_n^L$ [$\mu$m]</th>
<th>$L_n^S$ [$\mu$m]</th>
<th>$L_v$ [$\mu$m]</th>
<th>$L_{Te}$ [$\mu$m]</th>
<th>$L_{Ti}$ [$\mu$m]</th>
<th>$L_e$ [eV]</th>
<th>$T_e$ [eV]</th>
<th>$L_n$ [$\mu$m]</th>
</tr>
</thead>
<tbody>
<tr>
<td>Silicon</td>
<td>0</td>
<td>0.33</td>
<td>0.17</td>
<td>18.</td>
<td>37.5-70</td>
<td>40.30</td>
<td>0.25</td>
<td>0.06</td>
<td></td>
</tr>
<tr>
<td></td>
<td>260</td>
<td>7.</td>
<td>22.</td>
<td>7.5-10^6</td>
<td>16.5-596</td>
<td>15.6-16.7</td>
<td>—</td>
<td>—</td>
<td></td>
</tr>
<tr>
<td></td>
<td>510</td>
<td>13.</td>
<td>36.</td>
<td>35.</td>
<td>30-10.5</td>
<td>35-29.</td>
<td>—</td>
<td>—</td>
<td>12.5±3.</td>
</tr>
<tr>
<td></td>
<td>760</td>
<td>21.</td>
<td>49.</td>
<td>6.9-10^6</td>
<td>17.7-426</td>
<td>14.7-12.2</td>
<td>—</td>
<td>—</td>
<td>23±7.</td>
</tr>
<tr>
<td></td>
<td>1260</td>
<td>38.</td>
<td>65.</td>
<td>6.2-10^6</td>
<td>20.1-617</td>
<td>19.7-19.5</td>
<td>—</td>
<td>—</td>
<td>43±1.</td>
</tr>
</tbody>
</table>

Notes: $L_n^L$, $L_n^S$ are the electron density scalelength for planar and spherical plasma expansion. $L_v$ is the expansion velocity scalelength. $L_{Te}$ is the electron temperature scalelength. $L_{Ti}$ is the ion temperature scalelength for linear plasma expansion. Values in bold typeface are for the plasma parameters after 1 ps of heating by the laser interaction pulse.
E- Laser focus displacement

Some CCD camera images of the specularly reflected laser pulse on Silicon targets were taken for P-polarized single and double laser pulses. Figure III.26 shows a schematic of the measurement geometry, and under it, one image for a single pulse and two images for double pulses separated by $\Delta \tau = 1260$ ps. The single pulse image shows only a oval shape, the specular reflection of the single laser pulse at focus on the solid target surface.

![Schematic of measurement geometry]

**Fig. III.24 - CCD images of the incident laser focus reflected on a Silicon planar target for a single pulse and for a double pulse separated by $\Delta \tau = 1260$ ps.**

The target angle of incidence is $\theta = 55^\circ$. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The image insets for the single pulse (spp-07) and the second double pulse (dpp-08) are shown at twice the magnification than the larger pictures. The fainter spot in the double-pulse images corresponds to the preforming laser pulse. The darker spot overimposed on it is the interaction pulse reflection.
The double pulse image shows two superimposed oval shapes. The faintest one is the specular reflection of the preforming laser pulse focus on the Silicon solid target surface. The darkest oval is the specular reflection of the interaction laser pulse on the critical surface of the expanding Silicon plasma. The contrast ratio between the interaction and preforming pulse intensities is 3. The separation between the shapes is of $350 \pm 25 \, \mu m$. Using this measured distance between reflected images a direct estimation of the plasma's critical surface motion in the interval between the incidence of the preforming pulse and the incidence of the interaction pulse towards the observer can be made.

Assuming planar plasma expansion (with both angle of incidence and reflection supposed to be unchanged during the pulse duration), the Silicon plasma's critical surface is calculated to move towards the observer at a speed of

$$u_{\text{crit}} = \frac{d_{\Delta \tau} \sin \theta}{\Delta \tau} = \frac{350 \mu m \cdot \sin 55^\circ}{1260 \, \text{ps}} = 1.4 \pm 0.1 \cdot 10^7 \, \text{cm/s}.$$ 

This measurement is in close agreement to the LILAC planar expansion simulation that predict that the critical surface is moving towards the observer at $v_{\text{exp}} = 1.7 \cdot 10^7 \, \text{cm/s}.$

Using the electron temperature near the critical surface for $\Delta \tau = 1260 \, \text{ps}$ separated pulses, calculated using the integrating sphere measurements, therefore the sound speed near the critical surface is

$$c_s \approx c_{\text{crit}} = 9.8 \cdot 10^5 \sqrt{\frac{Z T_e}{A}} \, \text{[cm/s]} = 9.8 \cdot 10^5 \sqrt{\frac{14 \cdot 94 \, eV}{28}} \, \text{[cm/s]} = 6.7 \cdot 10^6 \, \text{cm/s}.$$ 

The LILAC predicted ion-acoustic speed for this case is $c_s = 7.0 \cdot 10^6 \, \text{cm/s}.$ The Silicon plasma's critical surface is moving at approximately $2.1 \cdot c_s$, after $1260 \, \text{ps}$ of plasma formation.

The planar LILAC simulation correlate better than the LILAC spherical simulations to the density scalelengths and electron temperatures measured using the integration sphere, and also to the propagation speed of the critical surface, measured by the image separation of specularly reflected light on this surface. All the comparisons between measured data and theoretical predictions described in chapters IV and V will be done using the corresponding values provided by the LILAC planar plasma expansion simulations.
REFERENCES FOR CHAPTER III


10 W. L. Kruer, ibid., ch. 5, pp.48-52.

11 W. L. Kruer, ibid., ch. 5, p. 54, eqn.5.26.


14 W. L. Kruer, ibid., ch. 4, pp.39-43.


39 Diode and acquisition electronics provided by Robert Boni.


42 This reflectivity was measured with a Perkin-Elmer spectrophotometer with light of wavelength in the spectral range 1030-1070 nm.


IV- Results and Analysis

In Chapter IV, the SBS spectral and energy measurements are shown and their features classified and compared with the SBS theory presented in Chapter II.

IV.1 - Silicon Targets

The targets are solid planar strips of one millimeter thick pure amorphous polished silicon (Si, A=28.1, Z=14), approximately one inch long and a 1/4 inch wide. The size of the target accommodates ~120 shots, with each laser shot crater separated by ~1 mm (twice the crater size). All SBS laser shots interacted with a undamaged fresh surface. Simultaneous measurements of the total backscattered energy and spectrum were done. Also, from the integration of the spectrum, the energy associated to the main spectral features was calculated.

A- Backscattered Reflectivity

The total energy of the backscattered light, $E_{bak}$, was measured with a large area UDT-100 PIN diode. A description of this diagnostic is given in section III.7.Ba and a schematic of it is in fig. III.14. The PIN diode gives its energy measurements in counts. The calibration of this device, shown in fig. III.19a, allows conversion of the instruments’ units of energy measurement to energies expressed in milliJoules (mJ). These measurements will be commonly referred in the following text as PIN energy or reflectivity measurements.

Simultaneously, the energies of the spectral features are calculated by the integration of the area under these features found in the spectra obtained with the 1-m Czerny-Turner spectrometer and detected by an EG&G/PARC Model 1471A Optical Multichannel Analyzer/Photo Diode (OMA/PD). A description of this diagnostic is given in section III.7.Bb and a schematic of it is in fig. III.14. The OMA/PD gives its energy measurements in nm-counts. The calibration of this device, shown in fig. III.19b, allows conversion of the instruments’ units of energy measurement to energies expressed in milliJoules (mJ). These measurements will be commonly referred in the following text as OMA energy or reflectivity measurements.
The backscattered intensities (expressed in W/cm²) are calculated by dividing the PIN and OMA measured energies by the laser pulse width (FWHM= 1.6 ps), and the effective target area (at 55° incidence and for 1/e intensity. See section III.5 and fig. III.15) is

$$A_{\text{target}} = (11 \pm 0.6) \times (17 \pm 1.0) \text{μm}^2/\cos 55° = (1.0 \pm 0.1) \cdot 10^{-5} \text{ cm}^2.$$

**a-Single laser pulses**

The single pulse total backscattered energy $E_{\text{bak}}$ (with associated intensity $I_{\text{bak}} = E_{\text{bak}}/A_{\text{target}} \tau_{\text{pulse}}$) and its corresponding reflectivity

$$\mathcal{R}^{\text{single}} = \frac{E_{\text{bak}}}{E_{\text{int}}} = \frac{I_{\text{bak}}}{I_{\text{int}}} \quad (\text{IV.1.1})$$

were measured for the interaction of single laser pulses with planar silicon targets. Figure IV.1 shows a typical plot of the total backscattered intensity $I_{\text{bak}}$ for both P-polarized (filled squares fitted with a full line) and S-polarized (diamonds fitted with a dashed line) interaction laser pulses of intensity $I_{\text{int}}$.

![Single laser pulse-Silicon](image)

**Fig. IV.1 - Total backscattered intensity—PIN diode measured—for single laser pulses interacting with a planar Si target.**

The target angle of incidence is 55°. The laser pulse has $\lambda_L = 1053.2$ nm, FWHM= 1.6 ps.
In both the P- and S- polarization states, the total backscattered intensities appear to grow strictly linearly with the increasing incident laser intensity, with a behavior similar to that of the no SBS region in figs. II.13-14. The backscattered reflectivities are constant and small,
\[ R_p^{\text{single}} = [5.0 \pm 0.1] \cdot 10^{-5}, \]
and
\[ R_s^{\text{single}} = [7.4 \pm 0.2] \cdot 10^{-5}. \]

The reflectivity for S-polarization, \( R_s^{\text{single}} \), is approximately 50% higher than the reflectivity for P-polarization, \( R_p^{\text{single}} \). This difference in reflectivities can be explained by the presence of resonance absorption when P-polarized pulses are used. As this measurement was done at an angle of incidence of 55°, the effect of resonant absorption in plasma formation (first) P-polarized laser pulses is maximum.¹

This measurement shows that the backscattered light resulting from the interaction of single laser pulses with planar silicon targets originates only from the Lambertian reflection of these pulses on the irregularities of the target surface, plus optical noise coming from reflection on the surfaces of the laser focusing and diagnostics light collection optics.

**b-Double pulses**

Introducing a second interaction laser pulse an interval \( \Delta \tau \) after the first preforming laser pulse shows a markedly differently behavior of the measured total backscattered intensity \( I_{b\text{ak}} \) and reflectivity \( R^{\text{double}} \) compared to the single pulse measurements, when the interaction intensity \( I_{\text{int}} \) is increased.

The measuring PIN diode cannot discriminate between the light scattered by the Silicon target originated with the preforming laser pulse, and the light scattered by the expanding plasma or Lambertian scattered by the critical density surface coming from the interaction laser pulse. Thus the measured reflectivity associated with the total backscattered energy for double incident laser pulses must be defined as

\[ R_{\text{PIN}}^{\text{double}} = E_{b\text{ak}}/(E_{\text{for}} + E_{\text{int}}) = I_{b\text{ak}}/(I_{\text{for}} + I_{\text{int}}). \]  \hspace{1cm} (IV.1.2)

However, when the backscattered energy is measured through integration of the area under features in the spectra taken with the OMA/spectrometer set-up (see section III.7.Bb), the reflectivity is defined as

\[ R_{\text{OMA}}^{\text{double}} = E_{b\text{ak}}/E_{\text{int}} = I_{b\text{ak}}/I_{\text{int}} \]  \hspace{1cm} (IV.1.3)
The reflectivities $R_{\text{PIN}}^{\text{double}}$, $R_{\text{OMA}}^{\text{double}}$, as defined in IV.1.2.3 not necessarily have the same value, even when measured simultaneously. $R_{\text{PIN}}^{\text{double}}$ gives the backscattered reflectivity of both the preforming and interaction pulses. $R_{\text{OMA}}^{\text{double}}$ represents the backscattered reflectivity of only the interaction laser pulse. But both reflectivities, even when different in value, experience qualitatively and quantitatively the same change with the appearance of SBS from the interaction pulse, therefore providing an useful tool to determine the value of the SBS threshold.

Figs. IV.2a-e show plots of the PIN and OMA measured backscattered reflectivities with respect to the interaction (P-polarized) laser pulse intensity for five different interval $\Delta t$ between formation and interaction laser pulses. The full diamonds represent the PIN measurements, the full circles represent the OMA measurements, and the $\times$-s represent the single pulse PIN measurements taken in the same experiment.

These plots show that the measured reflectivities are differentiated in two regions with increased interaction intensity. This behavior is related to the description of the measurement of the SBS backscattered signal described in section II.6B and shown in figs. II.13-14.

![Silicon—$\Delta t = 260$ ps](image)

**Fig. IV.2a - Reflectivity for double laser pulses separated by $\Delta t = 260$ ps interacting with a planar Si target.**
The target angle of incidence is $55^\circ$. The laser pulse has $\lambda_L = 1053.2$ nm, FWHM = 1.6 ps and it is P-polarized.
Fig. IV.2b - Total reflectivity for double laser pulses separated by $\Delta \tau = 510$ ps interacting with a planar Si target.

Fig. IV.2c - Total reflectivity for double laser pulses separated by $\Delta \tau = 760$ ps interacting with a planar Si target. The target angle of incidence is $55^\circ$. The laser pulse has $\lambda = 1053.2$ nm, FWHM = 1.6 ps and it is P-polarized.
Fig. IV.2d - Total reflectivity for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Si target.

Fig. IV.2e - Total reflectivity for double laser pulses separated by $\Delta \tau = 1260$ ps interacting with a planar Si target. The target angle of incidence is 55°. The laser pulse has $\lambda_L = 1053.2$ nm, FWHM = 1.6 ps and it is P-polarized.
It is clear that for the backscattered reflectivity $R_\text{double}$, two very different regions appear in each of the figs. IV.2a-e. For interaction laser pulse intensities $I_{\text{int}}$ below certain threshold value $I_{\text{thr}}$—value which varies with the interval $\Delta \tau$ between formation and interaction laser pulses—, the measured reflectivities $R_\text{double}$ are distributed approximately in a flat line of reflectivity value similar to those of the single laser pulse backscattered reflectivities. These data points are below the SBS threshold. This region shows no SBS, only Lambertian scattering from the target critical surface.

With increased interaction intensity, some data points start to appear with reflectivities distributed on a region well over the single-pulse reflectivity line. For interaction laser pulse intensities $I_{\text{int}}$ higher than the threshold intensity $I_{\text{thr}}$, the measured reflectivities $R_\text{double}$ experience a fast growth with increasing interaction laser pulse intensities $I_{\text{int}}$. This region over the threshold intensity $I_{\text{thres}}$ can be associated to linear SBS when correlated to the corresponding spectral measurements. There is no perceptible boundary leading to a differentiated region showing the non-linear (absolute) saturation of SBS.

**B- Spectra**

Simultaneously with the backscattered energy measurements described in the previous sections, the backscattered light spectrum—integrated over the interaction time of the pulse with the plasma—was obtained with the 1-m Czerny-Turner spectrometer and detected by an EG&G/PARC Model 1471A Optical Multichannel Analyzer/Photo Diode (OMA/PD) and converted and stored digitally with a GPIB/PC set-up. A description of this diagnostic is given in section III.7.Bb, with a schematic in fig. III.14.

At the beginning of every experimental run, using the same 50% mirror used for PIN and OMA/PD energy calibration (as seen in section III.7.C and fig. III.18), the spectra of the incident laser pulse was taken without any target present inside the Zeta chamber. An averaged spectrum—sharply clean, appearing as short dashed line in the following spectra—was then calculated and used afterwards as a fiducial for the spectra of single and double laser pulses interacting with targets in the Zeta chamber.

**a-Single laser pulses**

Prior to every set of double laser pulses, the spectra of single shot, P-polarized pulses at different incident laser intensities ($10^{14}$ W/cm$^2 \leq I_{\text{int}} \leq 5 \cdot 10^{16}$ W/cm$^2$) was obtained for later comparison between spectral measurement sets.
All the single laser pulse measured spectra have the same features (fig. IV.3). Compared to the position and shape of the incident laser fiducial, the resulting spectra has the following characteristics:

- The measured signal is low and very noisy, due to the low reflectivity (as seen in section IV.1Aa, $R_p^{\text{single}} = 5 \cdot 10^{-5}$) of the Lambertian scattering on the target plus optical noise coming from the diagnostics.

- The spectral peak is slightly blue-shifted ($< 0.5$ nm) with respect to the laser fiducial, resulting from the motion towards the observer (see fig. II.12) of the target surface causing the Lambertian reflection of the incident laser pulse.

\begin{figure}[h]
\centering
\includegraphics[width=0.5\textwidth]{spectrum.png}
\caption{Backscattered spectrum (averaged) for single laser pulses interacting with a planar Si target.}
\end{figure}

The target angle of incidence is 55°. The laser incident pulse has FWHM = 1.6 ps and is P-polarized. The bluer side of the spectrum is at the left of the plot. The amplitude of both the single pulse (full line) and laser fiducial (short dashed line) have been normalized to one to allow an easier comparison of both the position of peaks and their shapes.

- The backscattered Lambertian feature is slightly narrower ($\Delta \text{FWHM} < 1.2$ nm) than the laser fiducial.

No additional spectral features were observed besides the one described above, even at the highest incident laser intensity levels achievable with the TTT-CPA laser. Therefore the
conclusion is that no SBS was observed for P-polarized, single shot laser pulses interacting with planar Silicon targets.

**b-Double laser pulses with fixed target angle of incidence**

When a second—referred as the interaction—laser pulse is introduced after an interval Δτ from the first—plasma preforming—laser pulse, the general appearance of the time integrated spectra of the light backscattered by the planar silicon targets changes markedly in comparison to the single pulse spectral measurements, when the interaction intensity I_{int} of this second pulse is increased.

Several sets of time-integrated spectral measurements for different temporal separations Δτ between the preforming and interaction laser pulses were taken. These particular values of Δτ—260, 510, 760, 1010 and 1260 ps—were chosen to allow a wide range of plasma scalelength sizes for different expansion times between formation and interaction; also due to operational constraints in the pulse splitting, delaying and stacking set-up (see section III.3) these values allowed easier determination of the movable mirror position (see fig. III.11). In all cases P-polarized incident laser light was used.

Figures IV.4a-e show the typical changes in the spectral appearance of the backscattered light from a planar silicon target at 55° with respect to the normal when interacting with double laser pulses—separated by Δτ= 1010 ps—with increasing interaction pulse intensity I_{int} \(10^{14} \text{ W/cm}^2 \leq I_{int} \leq 5 \cdot 10^{16} \text{ W/cm}^2\). In all figures no attempt has been made to normalize the peak amplitudes (appearing in OMA counts) of the features shown.

In fig. IV.4a, the spectrum shows a single feature, resulting from the Lambertian reflection on the target critical surface plus the diagnostics' optical noise (< 10^7 W/cm^2) of both the preforming and interaction pulses. This feature has essentially the same characteristics as the ones described for single laser pulse interactions (see section IV.1Ba and fig. IV.3); from hereafter it will be called the Lambertian feature (L).
Fig. IV.4a - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude.

Fig. IV.4b - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude.
When the interaction pulse intensity $I_{int}$ is increased, keeping the preforming intensity $I_{for}$ at the same level, a second feature starts to appear to the blue side of the Lambertian feature (L), as seen in fig IV.4b. The threshold intensity level $I_{thres}$ (≈ $8.5 \times 10^{15}$ W/cm$^2$, in this case) for Brillouin backscattering has been barely crossed. This feature will be referred hereafter as the Brillouin backscattered one (B).

Increasing the interaction pulse intensity $I_{int}$ over the SBS threshold even further (but again keeping the preforming intensity $I_{for}$ at the same level of the previous shots), the Brillouin backscattered feature (B) appears sharper and well defined and of increased amplitude, similar in width and appearance to the Lambertian feature (L, as seen in fig IV.4c). The overall spectrum looks qualitatively similar to the upper right, light framed plot in fig. II.13.

This sharpness means that the Brillouin scattering region of resonance in the plasma is limited in size. A large region of SBS resonance appears in the spectrum as a wide, blunt feature. Each part of an extended region of resonance—each one characterized with different plasma density and temperature—adds its contribution of differently SBS and Doppler shifted light to the spectrum. The integration of all this dissimilar scattered light

![Diagram](image)

**Fig. IV.4c - Backscattered spectrum for double laser pulses separated by $\Delta \tau=1010$ ps interacting with a planar Si target.**

The target angle of incidence is 55°. The incident laser pulse has $\lambda_i=1054$ nm, FWHM=1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude.
must appear as widely dispersed, with the amplitude of the feature diminished due to the spectral dispersion of the scattered light intensity.

With increasing of interaction intensity $I_{\text{int}}$ over threshold (but always keeping the preforming intensity $I_{\text{off}}$ the same), a new feature—bluer, smaller in amplitude, sharp and well defined—is added to the other two (see fig IV.4d), similar in characteristics to both the Lambertian (L) and the Brillouin backscattered (B) features. As seen in section II.6A and fig II.12, this feature results from Brillouin sidescattering of the specular reflection of the incident laser light on the silicon target’s critical surface. Therefore, it is referred as the secondary Brillouin sidescattered one (S). The overall appearance of this spectrum, as seen in fig. IV.4d, is qualitatively similar to the lower right, heavy framed plot in fig. II.13.

Near the highest interaction intensity levels available with the TTT-CPA laser ($\sim 5.10^{15}$ W/cm$^2$), both the amplitude and the width of the Brillouin back (B) and sidescattered (S) features become very large, starting to merge with the Lambertian feature (L) (as seen in fig IV.5e), and spreading farther out the bluer side of the spectrometer window.

![Graph showing Laser and Scattered data with peaks L, B, and S](image)

**Fig. IV.4d - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Si target.** The target angle of incidence is 55°. The incident laser pulse has $\lambda = 1054$ nm, FWHM=1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak (L) of the scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.4e - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda_\text{L} = 1054$ nm, FWHM=1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak (L) of the scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).

Similar evolution of the backscattered features with increasing interaction intensity (at the same level of preforming intensity) can be found for the other measured spectra taken at different double pulse intervals $\Delta \tau$—260 ps (figs. IV.5a-d), 510 ps (figs. IV.6a,b), 760 ps (figs. IV.7a-d).
Fig. IV.5a,b - Backscattered spectrum for double laser pulses separated by \( \Delta t = 260 \) ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has \( \lambda = 1054 \text{ nm}, \) FWHM=1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).

- **OMA-07116021**
  - \( I_{\text{for}} = 2.5 \cdot 10^{15} \text{ W/cm}^2 \)
  - \( I_{\text{int}} = 8.8 \cdot 10^{15} \text{ W/cm}^2 \)

- **OMA-07116023**
  - \( I_{\text{for}} = 4.2 \cdot 10^{15} \text{ W/cm}^2 \)
  - \( I_{\text{int}} = 31 \cdot 10^{15} \text{ W/cm}^2 \)
Fig. IV.5c, d - Backscattered spectrum for double laser pulses separated by $\Delta t = 260$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM$= 1.6$ ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.6a, b - Backscattered spectrum for double laser pulses separated by $\Delta T = 510$ ps interacting with a planar Si target.

The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.7a, b - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 760$ ps interacting with a planar Si target. The target angle of incidence is $55^\circ$. The incident laser pulse has $\lambda_i = 1054$ nm, FWHM = 1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the corresponding scattered spectrum amplitude. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.7c, d - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 760$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda_l = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The laser spectrum is normalized to an amplitude of 500 counts, because the spectrum’s Lambertian feature has merged with the backscattered Brillouin feature, making itself indistinguishable. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
The back- and sidescattered Brillouin features do not appear every time that the intensity of the interaction pulse has crossed over the SBS threshold, due to the lack of control on the preformed plasma conditions for individual laser shots. With longer delay between pulses $\Delta \tau$, both the backscattered (B) and the sidescattered (S) Brillouin features and the Lambertian feature (L) appear more often together when the SBS threshold is crossed over, with corresponding increase of their spectral peak amplitude. Also for the higher interaction intensity levels, these features merge more frequently. For the longest interval measured, $\Delta \tau = 1260$ ps (see figs. IV.8a-c), the conditions in the preformed plasma after such a long expansion time are such that the SBS features are scattered at unpredictable wavelengths (but always at the bluer side of the Lambertian feature [L]). Sometimes there are more than two SBS features present in the measured spectra (see fig. IV.9d), even when other spectra taken for similar values of the preforming and interaction intensities do not have more than the usual one or two.

Further quantitative analysis (with comparison to the SBS theory developed in chapter II) of the measured SBS spectral features appears in section §IV.4B.

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**Fig. IV.8a - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1260$ ps interacting with a planar Si target.**

The target angle of incidence is 55°. The incident laser pulse has $\lambda_l = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The laser spectrum is normalized to the scattered spectrum's Lambertian peak. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.8b, c - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1260$ ps interacting with a planar Si target. The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the corresponding scattered spectrum. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
**Fig. IV.8d - Backscattered spectrum for double laser pulses separated by \( \Delta T = 1260 \text{ ps} \) interacting with a planar Si target.**  
The target angle of incidence is 55°. The incident laser pulse has \( \lambda_L = 1054 \text{ nm} \), FWHM=1.6 ps, and is P-polarized. The laser spectrum is normalized to the Lambertian peak of the corresponding scattered spectrum. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).

### IV.2 - Gold Targets

To compare the effects of a high-Z and ion charge\(^2\) species on SBS thresholds and spectra, similar measurements to those done for Silicon were done for Gold (Au, \( A = 197, \ Z = 79 \)) with an interval between laser pulses of \( \Delta T = 1010 \text{ ps} \).

The targets are solid planar strips of 2000 Å thick gold film deposited over a flat glass or silicon substrate, approximately one inch long and a 1/4 inch wide. The size of the target accommodates \(~120\) shots, with each laser shot crater separated by \(~1\) mm (twice the crater size). As for Silicon, all SBS laser shots interacted with a undamaged fresh surface. Simultaneous measurements of the total backscattered energy and spectrum were done. Also, from the integration of the spectrum, the energy associated to the main spectral features was calculated.

**A- Backscattered Reflectivities**

The single pulse total backscattered intensity \( I_{bak} \) and its corresponding reflectivity,
was measured for the interaction of single laser pulses with planar gold targets. Figure IV.9 shows a typical plot of the total backscattered intensity $I_{bak}$ for P-polarized (×s fitted with a short-dashed line) interaction laser pulses of intensity $I_{int}$. It can be seen that the (P-polarized) single pulse reflectivity for gold is in the same range of values as that for silicon.

![Gold-single pulse](image)

**Fig. IV.9 - Backscattered intensity for single laser pulses interacting with a planar Au target.**

The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.

Figure IV.10 shows a plot of the reflectivities associated with the measured PIN ($R_{\text{PIN}}^{\text{double}}$, defined in eqn. IV.1.2, and plotted in full diamonds) and OMA ($R_{\text{OMA}}^{\text{double}}$, defined in eqn. IV.1.3, and plotted in full circles) backscattered intensities with respect to the interaction laser pulse intensity for an interval $\Delta \tau = 1010 \pi \sigma$ between the plasma formation and interaction laser pulses. Also the reflectivities (×) taken with single pulse PIN measurements are plotted.

As in the silicon measurements, two very different regions appear in fig. IV.10. For interaction laser pulse intensities $I_{int}$ below the threshold value $I_{thr}$ — a value which varies with the interval $\Delta \tau$ between formation and interaction laser pulses —, the measured
reflectivities are distributed approximately in a line similar to that of the single laser pulse backscattered intensities, denoting the absence of any Brillouin scattering.

For interaction laser pulse intensities $I_{\text{int}}$ higher than the threshold intensity $I_{\text{thr}}$, the measured reflectivities $R_{\text{PIN}}^{\text{double}}$, $R_{\text{OMA}}^{\text{double}}$, experience a fast growth with increasing interaction laser pulse intensities $I_{\text{int}}$. This is an indication of the appearance of Brillouin scattering. The measured intensity threshold for Gold, $(\sim 16\pm0.1)\times10^{15}$ W/cm$^2$ is slightly higher than that the corresponding Silicon threshold, $(\sim 9\pm0.1)\times10^{15}$ W/cm$^2$, at $\Delta\tau = 1010$ ps, but as the initial plasma conditions might be different in each case, no conclusion can be ascertained.

![Gold-1010ps](image)

**Fig. IV.10 - Backscattered reflectivity for double laser pulses separated by $\Delta\tau = 1010$ ps interacting with a planar Au target.** The target angle of incidence is 55°. The incident laser pulse has $\lambda_l = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.

**B- Spectra**

As in the Silicon spectral measurements, prior to every set of double laser pulses, the spectra of single shot, P-polarized pulses at different incident laser intensities ($10^{14}$ W/cm$^2 \leq I_{\text{int}} \leq 5\times10^{16}$ W/cm$^2$) for later comparison between spectral measurement sets.
One single pulse spectrum sample is shown in fig. IV.11. All the single laser pulse measured spectra have the same features. Compared to the position and shape of the incident laser fiducial, the resulting spectra has the following characteristics:

- The measured signal is low and very noisy, due to the low reflectivity (as seen in section IV.2Aa, $R_p^{\text{single}} = 4.6 \cdot 10^{-5}$) of the Lambertian scattering on the Z-target plus optical noise coming from the diagnostics.

- The spectrum peak is slightly blue-shifted ($< 0.3$ nm) with respect to the laser fiducial, resulting from the motion towards the observer (see fig. II.11) of the target surface causing the Lambertian reflection of the incident laser pulse.

- The backscattered Lambertian feature has a similar width ($\Delta \text{FWHM} = 1.2$ nm) than the laser fiducial.

No additional spectral features were observed besides the one described above, even at the highest incident laser intensity levels achievable with the TTT-CPA laser. Therefore the conclusion is that no SBS was observed for P-polarized, single shot laser pulses interacting with planar Gold targets.

**Fig. IV.11 - Backscattered spectrum for single laser pulses interacting with a planar Au target.**

The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.
CHAPTER IV - RESULTS AND ANALYSIS

Introducing an interaction laser pulse an interval $\Delta \tau = 1010$ ps from the first plasma preforming laser pulse, changes the general appearance of the time integrated spectra of the light backscattered by the planar Gold targets markedly in comparison to the single pulse spectral measurements, when the interaction intensity $I_{int}$ of this second pulse is increased. In all cases P-polarized incident laser light was used.

Figures IV.12a-c shows the typical changes in the spectral appearance of the backscattered light from a planar Gold target at $55^\circ$ with respect to the normal when interacting with double laser pulses—separated by $\Delta \tau = 1010$ ps—with increasing interaction pulse intensity $I_{int}$ ($10^{14}$ W/cm² $\leq I_{int} \leq 5 \cdot 10^{16}$ W/cm²). In all of the figures no attempt has been made to normalize the peak amplitudes (appearing in OMA counts) of the features shown.

In fig. IV.12a, the threshold intensity level $I_{thres}$ (= $16 \cdot 10^{15}$ W/cm², in this case) for Brillouin backscattering has been barely crossed, so the Brillouin feature is small and broad.

![Graph showing backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Au target.](image)

**Fig. IV.12a - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar Au target.**

The target angle of incidence is $55^\circ$. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM=1.6 ps, and is P-polarized. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.12b,c - Backscattered spectrum for double laser pulses separated by $\Delta t = 1010$ ps interacting with a planar Au target.
The target angle of incidence is $55^\circ$. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM=1.6 ps, and is P-polarized. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).

Increasing the interaction pulse intensity $I_{int}$ over the SBS threshold even further (but again keeping the preforming intensity $I_{for}$ at the same level of the previous shots), the Brillouin backscattered feature (B) appears well defined and has an increasing amplitude.
but again wider than the Lambertian feature (L). The overall spectrum, as in fig. IV.12b, looks qualitatively similar to the upper right, light framed plot in fig. II.13.

At the highest interaction intensity levels available with the TTT-CPA laser (50·10^{15} W/cm^2), both the amplitude and the width of the Brillouin back (B) and sidescattered (S) features become very large, starting to merge with each other and with the Lambertian feature (L) (as seen in fig IV.12c), and spreading farther out to the bluer side of the spectrometer window range.

### IV.3 - PARYLENE-N (CH)

The effect of a multispecies plasma on SBS thresholds and spectra^3 were measured by using targets of 5000 Å thick Parylene-N (CH, A= 12+1, Z= 6+1) film, deposited over a flat glass or silicon substrate, with an interval between laser pulses of Δτ= 1010 ps.

The targets are approximately one inch long and a 1/4 inch wide. The size of the target accommodates ~120 shots, with each laser shot crater separated by ~1 mm (twice the crater size). As for Silicon, all SBS laser shots interacted with a undamaged fresh surface. Simultaneous measurements of the total backscattered energy and spectrum were done. Also, from the integration of the spectrum, the energy associated to the main spectral features was calculated.

#### A- Backscattered Reflectivities

The single pulse total backscattered energy E_{bak} and its corresponding reflectivity,

\[ R^{\text{single}} = \frac{E_{bak}}{E_{int}} = \frac{I_{bak}}{I_{int}} = (6.8 \pm 0.1) \cdot 10^{-4}, \]

was measured for the interaction of single laser pulses with planar Parylene-N targets. Figure IV.13 shows a typical plot of the total backscattered intensity I_{bak} for P-polarized (×s fitted with a short-dashed line) interaction laser pulses of intensity I_{int}. It can be seen that the (P-polarized) single pulse reflectivity for Parylene-N is one order of magnitude higher than the single laser pulse reflectivity values for Silicon and Gold. The CH surface film of Parylene-N targets was mechanically more unstable than the corresponding surfaces of Silicon and Gold targets. Therefore a less uniform plasma is expected to be generated. This lack of uniformity enhances the Lambertian scattering of the incident light, explaining the higher (P-polarized) single pulse reflectivity for Parylene-N (CH).
Fig. IV.13 - Backscattered intensity for single laser pulses interacting with a planar CH target.
The target angle of incidence is 55°. The incident laser pulse has λL = 1054 nm, FWHM =
1.6 ps, and is P-polarized.

Introducing double laser pulses for plasma preformation and interaction, fig. IV.14
shows a plot of the reflectivities associated with the measured PIN (R_{PIN}^{double}, defined in
eqn. IV.1.2, and plotted in full diamonds) and OMA (R_{OMA}^{double}, defined in eqn. IV.1.3,
and plotted in full circles) backscattered intensities with respect to the interaction laser pulse
intensity for an interval Δτ = 1010 ps between formation and interaction laser pulses. Also
the reflectivity level of the taken (P-polarized) single pulse PIN measurements is plotted
(short dashed line).

In contrast to the Silicon and Gold measurements, for Parylene-N (CH) it is not
possible to find any differentiation between a no SBS region and the start of the linear SBS
region in fig. IV.14. Therefore it is not possible to determine a value for the convective
SBS intensity threshold for Parylene-N.

Both the PIN and OMA signals and the associated reflectivity measurements also
include a high level of Lambertian-scattered related signal and reflectivity, due to the poor
optical quality of the Parylene-N film covering the targets. As it is seen in the following
section, the Lambertian feature (L) was superimposed on the Brillouin backscattered feature
(B) in all the Parylene-N spectra, so it was very difficult to determine from which of these
features the backscattered intensity was originating. Therefore the SBS appears to be above threshold for all of the measured double pulse intensities.

![Graph](image)

**Fig. IV.14 - Backscattered reflectivity for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar CH target.** The target angle of incidence is 55°. The incident laser pulse has $\lambda_{I} = 1054$ nm, FWHM=1.6 ps, and is P-polarized.

### B- Spectra

Prior to every set of double laser pulses, the spectra of single shot, P-polarized pulses at different incident laser intensities ($10^{14}$ W/cm$^2 \leq I_{\text{int}} \leq 5 \times 10^{16}$ W/cm$^2$) for later comparison between spectral measurement sets.

Similarly to the one seen in fig. IV.15, all the single laser pulse measured spectra have the same features. Compared to the position and shape of the incident laser fiducial, the resulting spectra is noisy, but the amplitude of the single interaction pulse spectra is usually higher than in the Silicon and Gold cases. The bad optical quality of the Parylene-N film enhances the reflectivity (as seen in section IV.3A, $R_{p}^{\text{single}} = 6.8 \times 10^{-4}$) of the Lambertian scattering on the planar solid target. The spectrum peak is slightly blue-shifted (< 0.5 nm) with respect to the laser wavelength. The backscattered Lambertian feature is wider ($\Delta \text{FWHM} \approx 1.2$ nm) than the laser fiducial.
Fig. IV.15 - Backscattered spectrum for single laser pulses interacting with a planar CH target.
The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.

No additional spectral features were observed besides the one described above, even at the highest incident laser intensity levels achievable with the TTT-CPA laser. Therefore the conclusion is that no SBS was observed for P-polarized, single shot laser pulses interacting with planar Parylene-N (CH) targets.

When introducing a second interaction pulse after the first preforming one, as in the Silicon and Gold case, the spectra shows differences compared to the spectra taken for single laser pulses, although these differences are harder to perceive. For all preforming and interaction intensities, the spectra taken shows some kind of additional feature, sometimes even to the red side of the ever-present Lambertian feature.

Figures 16a-d show some of the spectra taken for Parylene-N, with both backscattered and sidescattered Brillouin features. It was not possible to find an orderly progression of spectra similar to the ones in fig. IV.4. It is noticeable that the Brillouin backscattered (redder) feature is usually superimposed on the Lambertian feature.

This superposition of Lambertian (L) and Brillouin backscattered (B) features result in very large spectral intensity measurements, with the consequent associated highly reflectivities, as high as 0.5% (see fig. IV.14).
Fig. IV.16a, b - Backscattered spectrum for double laser pulses separated by $\Delta \tau = 1010$ ps interacting with a planar CH target. The target angle of incidence is 55°. The incident laser pulse has $\lambda_l = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
Fig. IV.16c, d - Backscattered spectrum for double laser pulses separated by $\Delta t = 1010$ ps interacting with a planar CH target.
The target angle of incidence is 55°. The incident laser pulse has $\lambda_l = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. The primary Brillouin backscattered peak (B) and the secondary Brillouin sidescattered peak appear to the blue side (left) of the Lambertian peak (L).
IV.4 - Common Features of the Experimental Data

Considering all the intensity and spectral data as a whole, some definitive qualitative and quantitative common patterns related to the Brillouin scattering instability appear.

A- SBS Intensity Thresholds

Collecting all the measured SBS thresholds for Silicon and plotting them versus the different delay between plasma formation and interaction pulses (see fig. IV.18), is clear that the laser interaction intensities needed to excite the SBS instability decreases with increasing plasma expansion time. This is expected because of the increase of the density and expansion velocity scalelengths, reducing the SBS convective threshold.

The theoretical SBS related thresholds for the Liu-Rosenbluth-White (LRW) isothermal (convective) and Short-Epperlein (SE) non-isothermal models are obtained by the application of the eqns. II.1.59-60 (for the dissipative threshold), II.2.28 (for the convective threshold), and II.3.11 (for the temporal threshold) to the plasma parameters calculated by LILAC simulations with and without plasma heating effects by the interaction laser pulse. These values are calculated for each LILAC cell in the simulated data considered as having a homogeneous plasma profile. Figure IV.17 shows typical plots of LILAC-calculated LRW SBS thresholds along the expansion profile of a plasma interacting with a laser pulse after $\Delta t = 510$ ps and $\Delta t = 1010$ ps of its formation. The minimal values in these profiles are taken for all considered intervals $\Delta t$ (260, 510, 760, 1010, and 1260 ps) as the corresponding SBS thresholds for that particular expansion time.

A comparison of the predicted convective LRW and SE SBS thresholds to the measured ones gives the following results,

- The LRW model applied to the non-heated LILAC simulation data (see $\times$s and long-dashed thick lines in fig. IV.18) correlates very well to the experimentally obtained SBS threshold data taken with the longest expansion times—760 ps, 1010 ps, 1260 ps. However the convective SBS thresholds are underestimated by a factor of about 1.5 to 2 for the shortest—260 ps, 510 ps—plasma expansion times.
**Fig. IV.17a** - LILAC-calculated LRW SBS thresholds for a Silicon plasma formed at $\Delta \tau = 510$ ps before interaction.

LRW is the Liu-Rosenbluth-White isothermal model. Full diamonds are for LILAC simulation data not heated by the interaction pulse. Full circles are for LILAC simulation data after 1 ps of heating by the interaction pulse.

**Fig. IV.17b** - LILAC-calculated LRW SBS thresholds for a Silicon plasma formed at $\Delta \tau = 1010$ ps before interaction.

LRW is the Liu-Rosenbluth-White isothermal model. Full diamonds are for LILAC simulation data not heated by the interaction pulse. Full circles are for LILAC simulation data after 1 ps of heating by the interaction pulse.
- The LRW model applied to the laser pulse heated LILAC simulation data (see +s data points joined by medium-dashed thin lines in fig.IV.18) is predicts SBS threshold within a factor of two over the measured SBS thresholds for all expansion times.

The best correlation between the measured SBS thresholds for Silicon is for the LRW model applied to non-heated LILAC simulation data; when the LRW model is applied to the heated (after 1 ps) LILAC simulation data, the predicted thresholds become slightly too high. This suggests that the SBS scattering in the preformed Silicon plasma undergoes before the first picosecond of interaction, while the plasma is cool enough to support the Brillouin instability. After that SBS would be switched off by the high plasma temperatures produced by self heating by the interaction pulse.

![Graph showing SBS threshold vs. delay between pulses for different scenarios.]

**Fig. IV.18 - Measured and calculated SBS thresholds for Silicon at different interval between laser pulses.**
The target angle of incidence is 55°. The incident laser pulse has $\lambda_i = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. LRW is the Liu-Rosenbluth-White isothermal model, SE is the Short-Epperlein non-isothermal correction.
- Applying the SE model to the heated LILAC simulation data for Silicon (see open circles connected by long dashed thick lines in fig.IV.18), it is found that the non-isothermal correction factor reduces the predicted SBS threshold by at least an order of magnitude below the measured Brillouin threshold.

The dependence of the convective SBS threshold on the plasma density scalelength is shown in fig. IV.19. Here the measured Silicon SBS thresholds are plotted versus the measured (via integration sphere; see § III.7D and fig. III.23) plasma density scalelengths (full squares). The predicted LILAC density scalelengths are also plotted (open circles for the LRW model applied to an unheated plasma; open diamonds for a plasma heated during 1 ps by the laser pulse). In all cases, the data points must follow the dependence

\[ I_{\text{SBS}} \propto 1/L_n. \]

Deviations on the value of the convective SBS threshold from this inverse dependence to the density scalelength can be explained by the presence of additional plasma expansion velocity \( L_v \) and species temperature \( L_{T_e}, L_{T_i} \) scalelengths (as shown in § II.2B and eqns II.2.27-28) that could not be measured and hence to be included in this plot.

![Graph showing SBS threshold intensity vs density scalelength](image)

**Fig. IV.19** - Measured and calculated (LRW model) SBS thresholds for Silicon at different electron density scalelengths.
The target angle of incidence is 55°. The incident laser pulse has \( \lambda_L = 1054 \text{ nm}, \) FWHM = 1.6 ps, and is P-polarized.
CHAPTER IV - RESULTS AND ANALYSIS

To see the effects of higher ion mass and degree of ionization, the Gold convective SBS intensity threshold measurements were obtained. Also similar measurements were performed for Parylene-N (CH) to study the behavior of a plasma consisting of two ion species of very different mass. These results, compared to the Silicon threshold data, are shown in Table IV.1, plus the theoretical predictions of the corresponding models used. It can be seen that the LRW model gets SBS convective threshold values closer to those measured in all cases—within a factor of two for Silicon and Gold.

Table IV.1 - Measured and calculated SBS thresholds (in units of $10^{15}$ W/cm²) for Silicon, Gold, and Parylene-N (CH) at $\Delta \tau = 1010$ ps.

<table>
<thead>
<tr>
<th></th>
<th>Silicon</th>
<th>Gold</th>
<th>Parylene-N</th>
</tr>
</thead>
<tbody>
<tr>
<td>Measured</td>
<td>8.5±2</td>
<td>16±2</td>
<td>&lt;6</td>
</tr>
<tr>
<td>LRW unheated</td>
<td>6.5</td>
<td>23</td>
<td>0.015*</td>
</tr>
<tr>
<td>LRW + 1ps heating</td>
<td>13.5</td>
<td>45</td>
<td>0.8*</td>
</tr>
<tr>
<td>SE corrected</td>
<td>0.2</td>
<td>0.37</td>
<td>—</td>
</tr>
</tbody>
</table>

Notes: The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM= 1.6 ps, and is P-polarized. LRW is the Liu-Rosenbluth-White isothermal model. SE is the Short Eperlein non-isothermal model.
* Values are calculated for the CH fastest growing mode.

To make a better determination of the convective SBS threshold, a more detailed knowledge of the thermal conditions in the preformed plasma must be previously obtained. Also the effects of interaction pulse-heating during the Brillouin scattering must be addressed, specially for the longer pulses used in Inertial Confinement Fusion.

Therefore it seems that the classic LRW isothermal model is a good predictor for the values of the SBS instability intensity thresholds, provided that the relevant thermal parameters (species temperature, density and speed profiles) of the plasma during interaction are known—experimentally or theoretically.

B- Spectral Features

Looking at the spectra taken at a particular delay between pulses $\Delta \tau$, it is found that the Brillouin backscattered and sidescattered features are localized primarily at the same wavelength from shot to shot. For Silicon and Gold, these spectral features are strongly shifted to the blue. The average measured spectral position of the Brillouin backscattered
features for Silicon is shown in fig. IV.20; the predicted spectral shifts associated to SBS, Doppler and ion-acoustic wake effects are also shown for both heated (thin) and unheated (medium-thick) LILAC simulation data for Silicon.

Usually, the strong blueshifting of the SBS features has been explained by the Doppler effect produced by the highly supersonic expansion of the plasma towards the observer (eqn. II.2.4 at § II.2A). Looking at fig. IV.20, it is clear that the motion towards the observer of the Brillouin resonant regions of the plasma does not produce a Doppler blueshift ($\Delta \lambda / \lambda_L = -v_{exp}/c$) high enough to compensate the Brillouin redshift ($\Delta \lambda / \lambda_L = 2c_s/c$), even for highly supersonic expansion speeds. The predicted spectral position for Brillouin scattered features undergoing Doppler blueshifting approximately superimposes on the incident laser wavelength. If the effect of the refraction index shift (see § II.2A) is considered, this additional term

$$\frac{\Delta \lambda}{\lambda_L} = -\frac{d}{c dt} \int_{z>z_{cr}} \sqrt{1 - \frac{n_e(z)}{n_{cr}} dz}, \quad (IV.3)$$

is comparable or even higher in magnitude than the corresponding Brillouin redshift, but reversed in sign. Therefore, a very strong blueshift is added to the Brillouin redshift term and to the Doppler blueshift term. As a result, the measured spectral positions of Brillouin backscattered features fall into the range of Brillouin plus Doppler plus refraction index shift values predicted using LILAC simulations for Silicon to describe the plasma expansion profile.

Figure IV.21 shows a plot of the spectral position of the peaks of the Brillouin back- and sidescattered features for a Silicon plasma interacting with a laser pulse after an interval $\Delta \tau = 1010$ ps past its formation. Even when there is a trend (a linear curve fit of these points gives $\Delta \lambda_{back}/\lambda_L = -0.00034 \cdot I_{int}[10^{15} W/cm^2]$, $R = 0.013$) for these peaks to blueshift further with increased interaction intensity, most of the backscattered Brillouin peaks remain positioned around 1050 nm. Similar behavior is found for the peaks of the secondary Brillouin sidescattered features (small solid diamonds).

Figure IV.22 shows a plot of the Brillouin back- and sidescattered plus Lambertian spectral positions for Silicon. These positions are separated approximately by the same wavelength interval, so it is expected that these features are originated in a region of resonance that it is at approximately the same distance of the target’s critical surface.
Fig. IV.20 - Measured and calculated absolute spectral shifts of Brillouin backscattered light for Silicon.

The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized. The measured shifts (filled dots) are averages taken over the entire set of spectral measurements for the particular delay between pulses. The error bars cover all the measured values Brillouin features spectral shift. The total shift results from the addition of SBS, Doppler and refraction index change contributions. Plot a is for LILAC unheated plasmas. Plot b is for LILAC heated at 1ps plasma data.
Fig. IV.21 - Measured spectral shifts of Brillouin back- and sidescattered light for Silicon at $\Delta \tau = 1010$ ps.
The target angle of incidence is 55°. The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.

Fig. IV.22 - Spectral shift of Brillouin backscattered and sidescattered features for Silicon.
The incident laser pulse has $\lambda_L = 1054$ nm, FWHM = 1.6 ps, and is P-polarized.
CHAPTER IV - RESULTS AND ANALYSIS

REFERENCES FOR CHAPTER IV


V. SUMMARY AND CONCLUSIONS

A SUMMARY OF the main features of the experimental data is presented and an analysis of the correlations between experimental results and theoretical predictions is performed. The validity of the theoretical models and their features are discussed in terms of the observed SBS data.

V.1 - SUMMARY

A series of experiments have been performed on the short pulse TTT-CPA laser,¹,² to study the laser intensity threshold of the Stimulated Brillouin Scattering (SBS) instability and its dependence on the plasma's density scalelength \( L_n \) and target composition. The experiments were performed at 1054 nm with intensities up to \( 5 \times 10^{16} \) W/cm². The backscattered and specularly reflected light resulting from the interaction of two high power picosecond duration laser pulses with solid Si, Au and CH planar targets was measured spectrally. The first pulse created a preformed plasma of density scalelength 30-300 laser wavelengths with which the second pulse interacted.

An ultrashort (1.6 ps) preforming laser pulse created a short scale length plasma. After a variable time delay (0-1300 ps), a second ultrashort (1.6 ps) interaction laser pulse excited the SBS instability in the quasi-isothermally expanded plasma. A critical surface was present in the plasma, which allowed a study of the behavior of SBS for near critical density (\( 0.5 \leq n_e/n_{cr} \leq 0.8 \)). The SBS reflectivity and spectral features were measured for several laser intensities, time delays and contrasts between preforming and interaction laser pulses.

V.2 - THE BRILLOUIN THRESHOLDS

From the SBS backscattered energy measurements for Silicon, Gold, and Parylene-N (CH) shown in chapter IV, the following conclusions can be found:

- Single short (~1 ps) laser pulses interacting with the very short scalelength (<1 \( \mu m \)) plasmas formed from solid Silicon, Gold and Parylene-N (CH) targets do not show the appearance of any Brillouin scattered signal in the range of incident laser
intensities \(10^{14} < I_L < 5 \cdot 10^{16} \text{ W/cm}^2\) used during this experiment (see § IV.1Aa-Ba).

- Short laser pulses interacting with moderate scalelength (>30 \(\mu m\)) expanded preformed Silicon, Gold and Parylene-N (CH) plasmas show a fast increase of the backreflected light intensity (see § IV.1Ab) when their intensity is over certain threshold value—which depends of the separation in time \(\Delta \tau\) between preformation and interaction laser pulses.

- The classical Liu-Rosenbluth-White (LRW)\(^3\text{–}^5\) isothermal model (see § II.2B) predicts (using LILAC simulated data) the measured Silicon and Gold SBS convective thresholds, within a factor of two (see § IV.4A and figs. IV.21-23). In order to get better agreement between the measured data and the predictions, a more comprehensive experimental knowledge of the plasma’s thermal state during the interaction time is needed.

- The values of the measured SBS intensity thresholds for Silicon are (within approximation) inversely proportional to the integration sphere measured (see § II.1B, § III.7D and fig. IV.23) plasma’s density scalelengths,

\[
I_{\text{SBS}}^{\text{meas}} \propto \frac{1}{L_n} \propto \frac{1}{\Delta \tau}.
\]

Deviations in the measured SBS threshold value from this dependence can be explained by the simultaneous presence of unmeasurable velocity and temperature gradients in the plasma.

- For Silicon, the Short-Epperlein (SE) non-isothermal model\(^6\) for SBS (see § II.1D) overestimates the correction—by one order of magnitude or more (see § IV.21)—to the interaction pulse-heated LRW’s SBS convective threshold. The SE model overestimates the correction—in one to three orders of magnitude—for the non-heated LRW’s SBS convective threshold for Silicon.

For Gold, the SE model correction again does not get close to the measured SBS convective threshold at \(\Delta \tau = 1010\) ps. Even when the LRW model calculated SBS threshold is bigger (by a factor of two) than the measured value, as in the Silicon case, it is the closest prediction to the measured value.

- For Parylene-N (CH), the predicted Brillouin threshold at \(\Delta \tau = 1010\) ps, using a modified LRW isothermal theory\(^7\text{–}^8\) for multispecies plasma (see § II.1A) is well
下面的测量范围为背散射的能量。对于所有测量与 Parylene-N (CH)，即使在最低可实现的相互作用束强度下，SBS 也观察到了。由于实验未能检测到 SBS 不稳定性，也不包括能量和光谱测量，因此 SBS 门槛无法在 Parylene-N (CH) 下找到。因此，关于多物种 LRW 模型的有效性，可以获得进一步信息。

V.3 - Brillouin Spectra

从 SBS 背散射光谱测量结果，对于 Si, Au, 和 Parylene-N (CH) 在章节 IV 中显示，可以得出以下结论：

- 单个短 (~1 ps) 激光脉冲与非常短的波长 (<1 μm) 形成的 Si, Au, 和 Parylene-N (CH) 等靶材的等离子不显示 Brillouin 散射的特征。在实验中使用的激光强度 (10^{14}-5 \times 10^{16} \text{ W/cm}^2) 中只有弱的、嘈杂的 Lambertian 特征 (近似激光波长 \lambda_L = 1054 \text{ nm}) 出现。

- 短激光脉冲与介于中等长度 (~25 μm) 扩张预形成的 Si, Au, 和 Parylene-N (CH) 等离子体只显示 Lambertian 散射的特征，当其强度低于某一强度阈值时 (参见 fig. IV.5a)。这些 Lambertian 特征不随频率宽度增加，在强度增加时，仅在幅度上线性增加。

- 短激光脉冲与介于中等长度 (~25 μm) 扩张预形成的 Si, Au, 和 Parylene-N (CH) 等离子体显示出局部的、蓝移的特征 (与 Lambertian 特征较近 \lambda_L = 1054 \text{ nm})，特征 (参见 § IV.1Bb 和 figs. IV.5b, IV.6-9) 当其强度超过某一阈值时，这取决于预形成时间 \Delta t 之间的间隔和相互作用激光脉冲的特征。这些特征可以识别为 Brillouin backscattering 的结果。

- 当相互作用激光脉冲强度进一步增加时，另一个局部的、甚至更多蓝移的特征出现在测量的光谱上 (参见 figs. IV.5c,e)。这些特征结果来自 Brillouin sidescattering，是未共振吸收的光特殊反射，以激光角度发生。
target's critical surface (see § II.6A). The relative spectral position of this sidescattered Brillouin feature with respect to the Brillouin backscattered one agrees with the results in § II.2 and § II.6.

Further increase of the interaction pulse intensity results in the spectral widening of all Brillouin features, their merging and even the appearance of additional ones.

- The spectral shift for the Lambertian features cannot be explained by the Doppler shift resulting from the expansion of the plasma critical surface towards the observer (see § III.7E).

- The SBS features show an extreme shift to the blue (see § IV.4B and fig. IV.24).

The Doppler blueshift caused by the plasma expansion towards the observer is not enough to compensate—even at highly supersonic speeds—the redshift inherent to any Brillouin scattering feature (see § II.2A). The additional blueshifting could result from the change experimented by the plasma's refraction index during the period between the arrival and departure of the interaction laser pulse into the SBS resonance region.

**V.4 - Conclusions**

SBS experiments using two ultrashort (~1.6 ps) laser pulses separated by a time interval between 100 and 1300 ps have some definite advantages over previous experiments using single or double long pulses (FWHM > 10 ns), as discussed in § I.2. Cleaner and more predictable density, temperature and expansion profiles can be obtained by the plasma formation pulse, while the second interaction pulse keeps the plasma illuminated long enough to reach convective saturation (see § II.3A). Some general conclusions can be reached:

- If the thermal profile of a plasma is well known at interaction time with a short laser pulse, the Liu-Rosenbluth-White isothermal model is able to predict the convective SBS intensity threshold.

Preforming a plasma with another short laser pulse provides a clean, quasi-isothermal plasma, of predictable characteristics. However, the use of the proper diagnostics to find not only the density, but also the expansion velocity and temperature profile of the plasma is needed to provide a better SBS threshold value prediction using the LRW isothermal model.
• The Doppler expansion—even at highly supersonic speeds—of the Brillouin resonant regions of the plasma cannot explain the high blueshifting of SBS features usually found in Brillouin backscattering experiments.

The additional high level of blueshifting present on the measured spectra could be explained by the change of the plasma’s refraction index during the time while the interaction laser pulse moves in and out of the SBS resonance region. Even using very approximate hydrodynamic models give a very reasonable prediction to the refraction index shift; although the use of more detailed, but computationally taxing kinetic models,\(^7\)\(^8\), coupled to a better knowledge of the plasma profile could provide a better description of the Brillouin scattered light spectrum.

• The measurement of Brillouin growth rates is very difficult due to the problems of controlling and reproducing the plasma profile at the interaction.

As a result is impossible to isolate and identify sets of data which have the same initial plasma profile at interaction time, so that only the interaction intensity could be the only variable in play when trying to find the corresponding growth rate. Very stable and reproducible laser pulse characteristics—as focus intensity profile and pulse peak intensity—, plus good quality target surface uniformity and homogeneity are preconditions to achieve a successful measurement of the SBS growth rates.
REFERENCES FOR CHAPTER V


APPENDIX

A- Measurements of Backscattered Light near 351 nm in Omega Long Scalelength Plasma Experiments

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Abstract

The backscattered light near the laser wavelength (351 nm) has been temporally and spectrally resolved during long scale-length plasma experiments on the 24 beam Omega laser system. CH plasmas with 0.5 mm density scale lengths and electron temperatures of order 1 KeV were formed with up to 22 beams in exploding foil targets. After the peak density decayed below critical, an interaction beam with variable delay was focused into the plasma through distributed phase plates at intensities up to $10^{15}$ W/cm$^2$. Smoothing by spectral dispersion was used in some shots. Under the above conditions $10^{-3}$ to $10^{-4}$ of the incident laser energy was backscattered through the lens and no evidence for Stimulated Brillouin Scattering (SBS) was observed. This was consistent with the calculated SBS velocity gradient threshold of $\sim 2 \times 10^{15}$ W/cm$^2$. 
A.1 - INTRODUCTION

The Stimulated Brillouin Scattering (SBS) instability has been extensively studied due to its potential as a serious energy loss mechanism in laser driven inertial confinement fusion plasmas. In this parametric instability, incident electromagnetic waves interact with the low frequency oscillation modes of the plasma, resulting in scattered electromagnetic waves—lower (but near) in frequency than the original—and ion-acoustic waves. Long-scale length plasmas (10^3-10^4 times the laser wavelength)—with smooth density gradients and consequently large regions of resonant coupling among waves—can lead to increased convective SBS gain and enhanced plasma reflectivity. SBS can reduce the amount of laser energy coupled to ICF pellets. Extensive theoretical and experimental work has studied the appearance of SBS in the infrared (1054 nm) and green (531 nm) wavelengths. For near ultraviolet (351 nm) wavelengths there are three relevant sets of results for long scale-length laser plasmas. R.P. Drake et al. detected enhanced reflectivity and frequency shifts compatible with the presence of SBS instabilities and have measured backscatter levels of 1% for green and UV at intensity levels of 4·10^{15} W/cm^2 on the NOVA laser system with both low-Z (CH— with plasma scale length L_n = 175 μm) and high-Z (gold — with L_n = 350 μm) solid targets. P.E. Young et al. irradiated CH targets with NOVA at similar intensity levels and measured 4% backscattered energy (Δλ = -5 to 0 Å decreasing with intensity, expansion speed = 7·10^7 cm/s, L_n = 350 μm). Goldman, Tanaka et al. (LLE-GDL) found Brillouin backscatter levels lower than 3% (Δλ = +3 Å, expansion speed = 0.5-1.5·10^7 cm/s) for UV at intensity levels of 10^{15} W/cm^2 with dextran foam targets (L_n = 500-1000 μm) and exploding CH foil targets (L_n = 50 μm) producing underdense plasmas. They also observed strong correlation between SBS and the two-plasmon decay instability.

The intensity threshold for SBS is lowest in a homogeneous plasma. SBS is a convective instability so gradients in the electron density, temperature and expansion velocity can raise the threshold intensity significantly. The regions where the scattered ion-acoustic and electromagnetic waves follow the Manley-Rowe relations become smaller. For the experiments described in this work, the laser intensity was above the density gradient instability threshold but below the velocity gradient threshold. As a consequence, we did not observe any strong evidence for SBS.

In this paper we report on the results of the measurement of backscattered light near the laser wavelength (351 nm) obtained on a series of long scale-length (L_n = 500 μm) plasma
experiments at LLE's Omega laser. No significant evidence for SBS was observed. In section A.2 a description of the experiment is given. Plasma formation with the Omega laser, the interaction beam characteristics and the diagnostics used for these experiments are discussed. In section A.3 the main features of the observed spectra and temporal measurements are shown. In section A.4 the results are discussed and analyzed. The homogeneous and inhomogeneous thresholds for the SBS instability, plasma reflectivity and Thomson background level are calculated for the typical parameters of these experiments. Finally, in section A.5, there is a summary of results and conclusions.

A.2 - EXPERIMENTAL CONDITIONS

The spectrum, energy and temporal history of backscattered laser light from an interaction beam was measured during long scale length plasma experiments. In this section, the formation and characteristics of the long scale length plasmas are reviewed and the experimental setup is described.

A- Long Scale Length Plasma Formation

The characteristics of the long scale-length plasmas have been described both theoretically and experimentally in Ref. 21 and are briefly reviewed here. The plasmas were preformed by the explosion of mass-limited, 6 μm thick CH (plastic) foil targets with a subset of eight primary beams (wavelength 351 nm, energy ≤ 60 J/beam, spot at target surface ≈ 450 μm FWHM). The targets had a 500 Å Al coating to ensure that the plasma formation commenced on the surface. As the plasma expanded, another subset of eight secondary beams, incident 0.6 ns after the primaries (with parameters identical to those of the primaries) and a subset of six tertiary beams (wavelength 1054 nm, energy ≤ 100 J/beam), peaking 1 ns after the primaries, were used to heat the plasma and maintain its temperature. The Gaussian pulse duration of each beam was 600-650 ps. The primary and secondary beams were outfitted with Distributed Phase Plates (DPPs) in front of the focusing lens (f = 60 cm) and defocused 1.65 mm past the target surface to ensure on-target illumination uniformity. With this set-up of the Omega laser a hot nearly isothermal plasma with electron temperatures of around 1 KeV. and density scale-lengths of the order of 0.5 mm was maintained for 1-2 ns.21
B- Interaction Beam

The interaction beam (wavelength 351 nm, energy ≤ 45 J/beam) was focused into the preformed plasma along the axis of symmetry normal to the original plane of the target. It was fired with time delays between 0.5 to 1.7 ns with respect to the peak intensity of the plasma-forming primary beams. With these interaction beam delay settings, there is enough time for the long-scale length, underdense plasma to evolve. The Gaussian pulse duration was 600 ps with a peak intensity of between $5 \times 10^{13}$ up to $10^{15}$ W/cm$^2$. This beam was always focused on the target through a distributed phase plate (DPP)$^{22,23}$—in some cases with cell sizes twice those of the heating beams—centered at the surface of the target with a spot size of 82 μm (for the smaller cell DPP). In addition beam smoothing by spectral dispersion (SSD)$^{24}$ was used on some shots.

C- Experimental set-up

The energy, spectrum, and temporal history of light backscattered through interaction beam focusing lens was measured. A glass pick-off (used for calorimetric measures of the beam) split 4% of the backscattered light coming from the Omega target chamber and sent into an aluminum off-axis spherical mirror ($f=160$ cm) which reflected it into the diagnostics. Either the reflected energy and spectra or the temporal history were recorded. The extremely narrow spectral features made it unnecessary to temporally resolve the spectrum.

For spectral and energy measurements this light was split with a second glass blank. The diagnostic arrangement is shown in Fig. A.1. Ninety-six percent of the light was injected through a glass diffuser into a 1 m Czerny-Turner spectrometer (Bausch & Lomb flat grating with 2400 groves/mm, 2400 Å blaze wavelength) using Aerographic 4421 UV film magazines. The absolute spectral line position and resolution were obtained by on each shot by including Ar$^+$ laser lines (at 351.1 and 351.4 nm) in the center of the film image. The remaining 4% of the light was measured in a large area filtered (three 10% plus one UV passband filters) and baffled UDT 10D PIN photo-diode biased at 45 V. The diode signal was calibrated by measuring the reflection off of an uncoated glass wedge followed by an absorber (calorimeter) in the interaction beam.

For temporal measurements the spectrometer was replaced with an unfiltered Photochron II type streak-camera, using an S-20 photocathode. The streak duration was 2
ns and the time resolution was better than 10 ps/line. The streak signal was recorded on Kodak Tri-X film developed at ASA 5000.

Both spectral and temporal images were digitized in a Perkin-Elmer PDS Model 1010 GMS microdensitometer. The digitized images were intensity converted with the PV-WAVE image processing language.

![Diagram of diagnostic setup](image)

Fig. A.1 - Diagnostics set-up for Omega Stimulated Brillouin backscattering experiments (E6696).

**A.3 - CHARACTERISTICS OF THE BACKSCATTERED LIGHT**

Figure A.2a-c shows backscattered intensity spectra from 3 typical interaction shots with a preformed plasma on a linear scale (upper plots) and semilogarithmic scale (lower plots). The spectra are dominated by reflection near the interaction beam wavelength. The maximum shift observed in this feature is less than 0.05
Fig. A.2a - Backscattered spectrum for OMEGA shot #23978. The backscattered spectra taken for long scalelength plasma interactions show little structure (E6639).
Fig. A.2b - Backscattered spectrum for OMEGA shot #24151. The backscattered spectra taken for long scalelength plasma interactions show little structure (E6638).
**Fig. A.2c - Backscattered spectrum for OMEGA shot #24258.**

The backscattered spectra taken for long scalelength plasma interactions show little structure (E6637).
nm. This is in contrast with spectrum shown in Fig. A.3 observed during the interaction with a solid gold sphere. In this case the central feature is blue-shifted by more than 0.05 nm and strong red-shifted features are observed up to 0.4 nm from the laser wavelength. The primary spectral feature shown in fig. A.2 is slightly wider than the incident laser spectrum. All of the spectral data show a dominant main peak close to the laser frequency (in no case this peak drifts more of 0.05 nm from the laser frequency). The average width of the primary feature for 56 shots was $0.010 \pm 0.001$ nm compared to $0.004 \pm 0.001$ nm width of the incident laser pulse. The reproducibility of the spectrum allow us to identify two additional redshifted features $\lambda_{r1} = 351+ 0.076 \pm 0.013$ nm, $0.015 \pm 0.002$ nm FWHM; $\lambda_{r2} = 351+ 0.017 \pm 0.013$ nm, $0.015 \pm 0.001$ nm FWHM). The intensities of these features, relative to the primary one, are shown in fig. A.2. For the highest intensities of the interaction beam one or more flat ($0.036 \pm 0.005$ nm FWHM) Gaussian-like peaks appear overlapped on the blue side ($\lambda_b = 351- 0.070 \pm 0.027$ nm) of the main peak. These features could indicate the onset of SBS instability for these plasmas.

Figure A.4 shows the fraction of the interaction beam energy backscattered through the lens as a function of the interaction beam energy. $10^{-3}$–$10^{-4}$ of the interaction beam energy was backscattered by flat targets under all conditions.

Figure A.5 shows a typical temporal history of the backscattered light near 351 nm taken with the streak camera. The width of the 600 ps interaction pulse is shown for comparison. There was no fiducial so the peak is centered on the signal. The backscattered light shows a broad, slightly asymmetric structure with less light reflected towards the end of the pulse.

**A.4 - ANALYSIS AND DISCUSSION**

The characteristics of the backscattered light can be summarized as follows. Between $10^{-3}$ and $10^{-4}$ of the interaction beam energy is backscattered through the lens. The backscattered spectrum is dominated by light within 0.05 nm of the laser wavelength and no rapid temporal structure was observed. These characteristics were the same whether SSD was used or not.
Fig. A.3 - Backscattered spectrum for OMEGA shot #24110 for the interaction with a Solid alignment target. Strong SBS was observed in solid target interactions (E6636).
Fig. A.4 - Fractional backreflected energy.
Approximately 0.1% of the interaction beam energy was reflected back and collected by the focusing lens (E6601).
Fig. A.5 - Temporal profile of OMEGA shots 24665, 24668. The streak camera integrated images of the backscattered light show a pulse duration comparable to the interaction beam pulse (E6602).
A- Total backscattered energy and reflectivity

The density and flow profiles calculated with SAGE—a 2-D plasma hydrocode—are shown in fig. A.7. The density profile can be approximated with a square shape. The reflectivity calculated\(^{26}\) from the gradient of the index of refraction of a square profile for a plasma density 10% critical is 1.8·10\(^{-3}\). This is consistent with the measured data (fig. A.4). For a smooth parabolic density profile, the reflectivity would be negligible as it is the spatial gradient of the index of refraction which determines the reflectivity. Thus it appears that the dominant feature in the backscattered spectrum results only from light which is reflected off of the density jump generated in the expanding plasma—as shown in the SAGE plasma profile. This small amount of backscattered energy indicates the absence of enhanced reflectivity due to SBS instabilities.

For gold spherical targets the observed reflectivity is also about 0.1%—due to the large absorption near the critical surface in the center of the targets—and the main features are one or two wide redshifted peaks which comprise most of the reflected energy. These signatures suggest the appearance of SBS in these targets, with the SBS signal possibly seeded by the backreflected radiation coming from the critical surface.

B- Thomson ion feature peak

The Thomson backscattering cross section was calculated from the equation 11.6.18 of Krall and Trivelpiece\(^{27}\) for a multispecies CH plasma.\(^{28}\) The total backscattered solid angle integration was performed on the cone subtended by the focusing lens of the target chamber and centered on the beam path and the result was integrated over all frequencies. The backscattered power is:

\[
P_{\text{Thom}} = P_{\text{inc}} \frac{2\pi (1 - \cos \alpha) r_{0}^{2}}{N_{e}} \int_{0}^{\infty} \frac{1 + \xi_{e} Z(\xi_{e})}{K^{2} \lambda_{D e}^{2} \sigma(K, \Delta \omega)} \sum_{\text{species}} \frac{N_{s}}{N_{e}} \frac{e^{-\xi_{e}^{2}}}{\sqrt{2\pi} K u_{s}} d\omega,
\]

where \(\alpha = 9.5\) degrees is the angle subtended from the axis of symmetry going from the center of the tank to the center of the lens, \(N_{e} = 9\cdot10^{19}\) electrons/cm\(^{2}\) is the number of electrons per unit area along the plasma length, \(\lambda_{Ds}\) is the Debye length for the plasma species \(s\) and \(r_{0} = 2.8\cdot10^{-13}\) cm is the classical radius of the electron.
Fig. A.6. - Thermal Thomson scattering ion peak energies.
The low-intensity features in the spectra have energies close to those expected from thermal Thomson scattering (E6641).
The dielectric function\(^{29}\)

\[
\mathcal{D}(K, \Delta \omega) = 1 + \sum_{\text{species}} \frac{1}{K^2 \lambda_p^2} \left[ 1 + \xi_s Z(\xi_s) \right]
\]

is written in terms of the plasma dispersion function \(Z(\xi_s)\), with

\[
\xi_s = \frac{\Delta \omega}{\sqrt{2 \pi} K u_s}.
\]

The energy in the small side peaks was calculated by comparing their integrals with data from the PIN diode energy measurements. These values are plotted in fig. A4, normalized to the value of the theoretical Thomson scattering intensity.

The reflected energy in the Thomson scattered wavelength peak placed at the blue side of the laser reflection peak has the same order of magnitude to the calculated Thomson scattering reflectivity. The position of this peak is blueshifted by \(\Delta \lambda = 4.5 \, \text{Å}\) from the stationary plasma value. This implies an expansion velocity of \(4 \cdot 10^7 \, \text{cm/s}\), which is consistent to the SAGE predictions. However, the corresponding Thomson scattered wavelength peak placed at the red side of the laser reflection peak was not always observed (see fig. A.1).

**C- SBS Threshold**

The most straightforward SBS threshold to calculate is that corresponding to a homogeneous and static density profile. From ref. 1 (eqn. 8.18), the expression for the minimum intensity needed for triggering the instability is

\[
\left( \frac{\nu_m}{u_c} \right)^2 \geq \frac{8}{k_\parallel L},
\]

resulting in the threshold intensity for homogeneous density plasmas

\[
I_{\text{min-hom}} \approx \left( \frac{m_e c u_c}{e} \right)^2 \frac{1}{\sqrt{1 - \frac{n_e}{n_c}}} \frac{\omega_0}{L},
\]
in which \( \omega_0 \) is the angular frequency of the incident EM wave, \( k_0 \) is the respective wavenumber, \( v_{os} \) is the electron quiver velocity, \( u_e \) is the electron thermal velocity, \( n_e \) the electron number density, \( n_c \) the critical density of the plasma for the laser frequency and \( L \) is the length of the plasma.

Exploding foil target plasmas have density, temperature and expansion velocity gradients along the normal to the foil’s surface, which limit the region of resonant coupling among the incident EM, scattered EM and ion waves. The threshold conditions are modified and new expressions appear for the minimum intensity. For a plasma with a density gradient, the homogeneous plasma length \( L \) is substituted by the density gradient scale-length \( L_n = n / (\partial n / \partial x) \) (Ref. 1, eqn. 8.19), which leads to

\[
I_{\text{min-deg}} = \left( \frac{m_e c u_e}{e} \right)^2 \frac{1}{1 - \frac{n_e}{n_c}} \frac{\omega_0}{L_n}.
\]

For a freely-expanding plasma, again the homogeneous length \( L \) is substituted, now by the velocity scale-length \( L_v = c_s / (\partial v_{ex} / \partial x) \) and from ref. 1 (eqn. 8.20) then the convective threshold intensity is

\[
I_{\text{min-vc}} = \left( \frac{m_e c u_e}{e} \right)^2 \frac{1}{1 - \frac{n_e}{n_c}} \frac{\omega_0}{c_s} \frac{\partial v_{ex}}{\partial x}.
\]

The three values obtained for \( I_{\text{min}} \) in different plasma regimes were calculated for an electron temperature in the 1 KeV range, 10% critical plasma density and 351 nm incident laser wavelength. This plasma parameters lead to ion sound speed \( c_s = 1.8 \times 10^7 \) cm/s, Landau damping rate \( \gamma_L = 2.1 \times 10^{12} \) s\(^{-1}\), and collisional rate \( v_{ci} = 1.9 \times 10^{11} \) s\(^{-1}\). These lead to a convective, velocity gradient threshold of 2 \times 10^{15} \) W/cm\(^2\) in these plasmas. The expansion velocities and density profiles have been obtained from SAGE calculations (see fig. A7).

The experiment used interaction beams less than 1.2 \times 10^{15} \) W/cm\(^2\), below the threshold intensity for the convective inhomogeneous velocity profiles predicted from the SAGE 2-D simulations (see fig. A7). Thus the lack of observed SBS is consistent with the inhomogeneous plasma thresholds.
Fig. A.7 - Calculated SAGE plasma density and expansion velocity profiles. The expansion velocity gradient was estimated using hydrodynamic profiles calculated by the 2-D SAGE code (E6640).
A.5 - CONCLUSIONS

We have measured the backscattered light from the interaction of a single 351nm laser beam with a preformed, subcritical plasma.

Even when the broad blue and red features appearing in spectra taken at higher intensity levels could suggest the onset of SBS instabilities there is no clear evidence of this instability in the frequency or time domain. The backscattered pulse duration is comparable to that of the incident pulse (600 ns) and approximately 0.1% of the interaction beam energy is backscattered through the focusing lens. This value correlates very well with those calculated for purely specular reflection from the nearly square plasma profiles derived from SAGE simulations, showing no additional sources of backreflected light, like SBS. At intensities near the SBS threshold, the appearance of broad blue- and redshifted features suggests the onset of the SBS instability, though this is not apparent in the backscattered energy or temporal history.

Several reasons can be found to explain the absence of SBS signals. The appearance of multiple modes in CH for ion-acoustic phonons increases the intensity threshold for SBS due to greater collisional and Landau damping effects\(^{25}\). The laser intensities used (\(\leq 1.2 \times 10^{15}\) W/cm\(^2\)) were below the convective velocity gradient threshold power, so the growth of the instability was constrained to too small a region of the expanding plasma. The use of distributed (random) phase plates (DPPs) in the interaction beam optical path homogenized the beam cross-section, impeding the appearance of filamentation or “hot spots” which could produce localized SBS. There are no noticeable effects of SSD (Smoothing by Spectral Dispersion) in the data, due to the mismatch between the larger SSD frequencies (3 GHz) to those of the ion-acoustic scattered waves.

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REFERENCES FOR APPENDIX A


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